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CHARM — A SIGNAL OF THE CONFINEMENT PHASE TRANSITION *

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ABSTRACT

I review some features of proposed signals of the quark-gluon plasma in order to motivate a new one: measuring the transverse momentum of charmed particles. A sharp reduction in the mean $p_\perp$ will show that the matter produced in a collision has cooled through a confinement phase transition. I argue that other mechanisms will not duplicate the effect.

INTRODUCTION

In this talk I will propose to you a new signal[1] of the quark-gluon plasma. More precisely, as you will soon see, it is a signal of the confinement phase transition itself, rather than of the QCD plasma which lies above it. In a nutshell, I suggest measuring the transverse momenta of charmed particles which emerge from the central collisions of heavy ions. My contention is that passage through the phase transition will damp these momenta dramatically. I leave it to the experimentalists here to judge whether such a signal is a simple one to detect; I leave it to all of you to judge whether this signal is any better than its predecessors, although I will argue that this is the case. I feel fairly secure, however, in claiming that this signal is new on the heavy ion scene. Even if the reconstruction of charmed particles proves impracticable, I hope that some of the theoretical ideas involved in the signal, in particular the focus on the dynamics of the phase transition, will carry over into new efforts to invent diagnostic tools.

I believe there to be a widespread dissatisfaction with the signals so far proposed for the quark-gluon plasma. In order to show the motivation for the charm signal, I would like to call your attention to some features which most of these signals have in common. Thus I will begin with brief and extremely biased discussions of four much-discussed diagnostic tools:

(1) Dileptons, meaning continuum $\mu^+\mu^-$ and e^+e^- pairs,

(2) J/ψ suppression,

(3) Strange particle production, and

(4) $\pi\pi$ correlations.

Just as the discussions are incomplete and unfair, so are the lists of references. I apologize to anyone who feels left out.

I will then proceed to my new (improved?) signal,

(5) Charm.

As is usual in theoretical discussions, I will have in mind the central rapidity region in the collision of very large nuclei at RHIC energies. This assumption, essentially one of simplicity, will still leave us with many details to study. Whether the physics to be discussed is useful for the current experiments is a question for the future.

CONTINUUM DILEPTONS

Dilepton emission from a quark-gluon plasma[2] is calculated via the simplest kinetic theory. The rate of emission is obtained by folding the cross section for $q\bar{q}$ annihilation into leptons with the corresponding flux in a gas of free quarks and antiquarks in thermodynamic equilibrium:

$$R^q(q\bar{q} \to l^+l^-) = \int \frac{d^3p_1}{(2\pi)^3}\, f(p_1) \int \frac{d^3p_2}{(2\pi)^3}\, f(p_2)\, \sigma(q\bar{q} \to l^+l^-; p_1, p_2)\, v_{\rm rel} \ , \tag{1}$$

where

$$f(p) = \left[e^{\beta p_\mu u^\mu} + 1 \right]^{-1} \tag{2}$$

is the Fermi-Dirac distribution in the moving fluid. The cross-section σ comes from the simple QED diagram for the annihilation of a free $q\bar{q}$ pair, since the central assumption of kinetic theory is that particles propagate freely between collisions. One takes this rate and folds it with the temperature profile which emerges from hydrodynamic calculations, also allowing for the gradual disappearance of plasma in the mixed phase. The pion gas into which the plasma is converted during the phase transition also emits dilepton pairs, and one calculates the rate $R^\pi(\pi\pi \to l^+l^-)$ via a similar formula, using vector dominance for an estimate of the cross section. Numerical results based on the Bjorken hydrodynamic expansion,[3] assuming a reasonable initial temperature of 250 MeV, show that (1) emission from the pion gas is greater than emission from the plasma when the pair mass is less than 2 GeV, and (2) the extrapolation of direct Drell-Yan production overshadows plasma emission for larger masses. Thus the simplest dilepton signal, namely the pair mass distribution, is beset by backgrounds.

Naturally, one can devise more complex signals to pick out the plasma contribution, such as[4] the lack of correlation with the beam direction (in contrast to Drell-Yan emission); also, one can look[2] at dilepton yields in conjunction with other data, such as total multiplicity, to derive more reliable signals. The background problems, however, are not my main objection. The point is Eq. (1), which is based on the assumption of a weakly coupled quark-gluon plasma. This assumption is based on QCD perturbation theory, which looks at first glance like the perturbation theory in finite-temperature QED. Unfortunately, it has been painfully clear for perhaps ten years that the resemblance is most superficial. I will expand on this point below.

J/ψ SUPPRESSION

The idea of J/ψ suppression by the plasma has been thoroughly described here by Sean Gavin, who may well have buried it. Let me add my two cents, with perhaps different emphases.

The original idea[5] was based on the observation that the potential between heavy quarks in the plasma should show the effects of Debye screening:

$$V_{c\bar{c}}(r) = -\frac{\alpha_{\rm eff}}{r}\, e^{-r/\lambda_D}, \tag{3}$$

where one-loop perturbation theory gives for the Debye length $\lambda_D = (gT)^{-1}$ (for a pure glue plasma). Note that, up to numerical factors, this is the same as the classical result $\lambda_D = (\beta e^2 \rho)^{-1/2}$ when the density of charges behaves as $\rho \sim T^3$, as it does in a gas of free massless gluons. Analysis of the Schrödinger equation in this potential (using, actually, lattice results for λ_D) showed that there is no bound state of a $c\bar{c}$ pair for $T > 230$ MeV. Thus J/ψ particles cannot survive in the plasma if it is hot enough. Later work[6] showed that the J/ψ yield should have strong $p_\perp$ dependence, since a $c\bar{c}$ pair created with large $p_\perp$ should escape the plasma before the constituents have time to separate beyond the J/ψ radius.

And in fact, up to a 50% suppression of J/ψ production has been seen in the CERN experiments,[7] with strong $p_\perp$ dependence.

Unfortunately, the signal may be faked by mechanisms that do not involve plasma creation.[8] As Sean has shown in detail, chemical reactions in the hadron gas such as $\pi + J/\psi \rightleftharpoons D + \bar{D}$ and $\pi + J/\psi \rightleftharpoons \eta_c + \pi$ can account for the overall suppression, and superposition of the effects of initial state interactions can account for the tilting of the $p_\perp$ distribution towards higher values. This picture casts doubt upon the uniqueness of the plasma explanation, and thus may destroy the value of the J/ψ yield as a plasma signal.

My quibble (and that of others) with this "conventional" explanation of suppression is that it requires a pion density of 1 fm^{-3} (in previous work[9] even 5 fm^{-3}) which is hardly that of a dilute gas of pions. It would not be unreasonable to call this "pion gas" a dilute plasma, and to consider the pion gas calculations to be merely a new approximation in dealing with the plasma: one that takes into account the correlations between quarks and antiquarks, and that estimates scattering off these $q\bar{q}$ pairs by using mesonic cross sections. Thus perhaps the alternative presented is still a plasma, albeit one with correlations beyond those expected in a weakly coupled system.

A truly conventional, non-plasma explanation of J/ψ suppression may emerge from the suggestion of Brodsky and Mueller,[10] who give arguments based on the parton model and analogies to QED. Here a $c\bar{c}$ pair are ripped apart by comoving partons, which are naturally denser in AA collisions than in pp. This picture even reproduces, if only qualitatively, the observed $p_\perp$ dependence. The distinction between this and other explanations is that no local thermal equilibrium is postulated. No plasma, no pion gas—just parton wave functions. Thus is the dramatic experimental signal relegated to the realm of the banal. Unless you like plain perturbative QCD, of course.

If a very dense pion gas points the way to an exceedingly complex quark plasma, so does, unfortunately, QCD perturbation theory[11] at $T \neq 0$. The screening length, so essential to the J/ψ suppression mechanism, offers the best glass through which to examine this can of worms. Consider a calculation of the gluon vacuum polarization tensor. The one-loop diagrams give a finite result which allows taking the external momentum to zero and thus calculating $\lambda_D^{-2} \equiv \Pi_{00}(q \to 0) = g^2T^2$ (in the quarkless, pure glue theory). Simple power-counting shows that an n-loop diagram gives a result of the form

$$\Pi_{00}^{(n)}(q) \sim g^4T^2 \left(\frac{g^2T}{q}\right)^{n-2}, \tag{4}$$

so that the $q \to 0$ limit poses a problem. In effect, the external momentum q is the only infrared cutoff in these diagrams, so they diverge when $q \to 0$. The apparent cure is to insert a gluon mass (a magnetic mass) into the propagators in the diagrams, so that a cutoff remains when $q = 0$. Now we have

$$\Pi_{00}^{(n)}(q) \sim g^4T^2 \left(\frac{g^2T}{m_{\rm mag}}\right)^{n-2}, \tag{5}$$

and all we need is the value of $m_{\rm mag}$. Unfortunately, calculation shows that $m_{\rm mag}$ is zero in one loop, and we can have at best $m_{\rm mag} = O(g^2T)$. Inserting this into (5) brings disaster—all the n-loop diagrams, for $n > 1$, give contributions of the same order in g. The Debye length is *not calculable* in perturbation theory.

The same problem appears, sooner or later, in the calculation of any quantity in perturbation theory. It is thus no exaggeration to say that, outside of some qualitative results from lattice calculations, the weakly coupled QCD plasma is a fiction.

This is a serious objection to many proposed signals of the quark-gluon plasma. Furthermore, note that this discussion applies even to extremely high temperatures where asymptotic freedom supposedly makes the effective coupling small. Who can say what happens around $T = 200$ MeV?

STRANGENESS

Strangeness enhancement,[12] at least in the central region, is not really a signal of the plasma at all. The reason is simple.[13] Consider a plasma in equilibrium at temperature T. The ratio of strange quarks to non-strange quarks and gluons is given approximately by

$$\frac{n_{s,\bar{s}}}{n_{q,\bar{q},g}} \simeq \frac{3}{14}\, e^{-m_s/T}, \tag{6}$$

where the prefactor comes from counting degrees of freedom and the Boltzmann factor contains the mass of the strange quark. If this plasma hadronizes quickly, the number of strange particles will equal the number of strange quarks, and the total number of particles will equal the total number of quarks and gluons because of entropy conservation. Thus Eq. (6) gives the ratio n_K/n_π directly in this case. On the other hand, a meson gas in equilibrium at the same temperature yields

$$\frac{n_K}{n_\pi} \simeq \frac{4}{3}\, e^{-(m_K - m_\pi)/T}, \tag{7}$$

which is *larger* for temperatures above 150 MeV. Either result is 3–4 times as large as the ratio measured in pp scattering, where thermodynamics is irrelevant. Thus enhanced strangeness production in the central region is a signal of rescattering and equilibration, but not of plasma creation.

More detailed studies,[13] attempting to discern the difference between the plasma and the hadron gas, have taken into account the departure from equilibrium in the hydrodynamic expansion. A typical formula for the divergence of the strange particle current (the strange particle creation rate) is

$$\partial_\mu n^\mu = \left(e^{-2\beta\mu} - 1\right) \int |\mathcal{M}|^2 f_1 f_2 f_3 f_4\; \delta^4(\textstyle\sum p_i), \tag{8}$$

where μ is the strange particle chemical potential which gauges the degree of departure from equilibrium. I have indicated schematically the matrix element $\mathcal{M}$ for creation and annihilation of $s\bar{s}$ pairs and the momentum distribution functions $f_i(p_i)$ (Bose-Einstein or Fermi-Dirac) for the species involved in each reaction. Naturally, this is nothing other than kinetic theory. Again, I call for caution in using such a formalism for the real QCD plasma.

$\pi\pi$ CORRELATIONS

Miklos Gyulassy has discussed pion interferometry here, so I will be exceedingly brief. As he has explained, one measures two-particle distributions and takes a Fourier transform in the relative momentum variable of the pairs. This should give information about the source of the pions: its size, its lifetime, and maybe even something more about its space-time structure. The great conceptual advantage of the method is that it does not depend on any assumptions about the dynamics of the source; all one needs is something that emits pions. Thus while pion interferometry will not probe the plasma itself, it might give indications of, say, bubble growth and droplet evaporation in the course of the phase transition.[14]

Unfortunately, pion interferometry has problems. Knowledge of source structure requires bump-hunting in the correlation function, and current experiments have not yielded enough statistics to allow this. Perhaps the higher multiplicities at RHIC will help. In addition, there are "background" problems which consist of contamination by late resonance decay, which Miklos has discussed at length.

Both the advantages and disadvantages of pion interferometry should be kept in mind when considering the discussion which follows.

CHARM

Charmed particles appear to be ideal candidates for diagnostic tools. The reasons are several:

- They can only be produced in the initial stage of the collision, when the nuclei are passing through each other and hard nucleon-nucleon collisions are taking place. This is because the mass of a charmed particle is much greater than any temperature which may reasonably be attained.[15]
- They move slowly, that is, at around or less than half the speed of light. This means that they do not escape the collision region early.
- Hence they experience the confinement phase transition, which is the part of the collision process that can perhaps best be described qualitatively, without detailed assumptions about the plasma.
- The motion of charmed particles may be characterized as Brownian motion,[16] since they are much heavier than the particles with which they are likely to collide. This gives an intuitively appealing picture for calculation.

In order to begin at the beginning, let us consider the motion of a charmed quark through a quark-gluon plasma. The quark experiences a drag force which damps its initial momentum, and at the same time diffuses in momentum space until its mean square momentum reaches the value appropriate for thermal equilibrium. Thus a quark which begins with a definite momentum $\vec{p}_0$ at $t = 0$ will at a later time have a momentum chosen from a distribution which satisfies

$$\langle \vec{p} \rangle = \vec{p}_0 e^{-\gamma t} \tag{9}$$

where the drag coefficient γ is given by (again from kinetic theory!)

$$\gamma = \lim_{\vec{p} \to 0} \frac{1}{p^2} \int d^3k \, w(\vec{p}, \vec{k}\,) \, \vec{p} \cdot \vec{k} \ . \tag{10}$$

I have denoted by w the transition rate for the heavy quark to go from momentum $\vec{p}$ to momentum $\vec{p} - \vec{k}$ by scattering off a plasma particle. The drag coefficient was calculated in Ref. 16 and it was found that $\gamma^{-1} \simeq 15$ fm at $T = 200$ MeV, if the calculation can be trusted at such a low temperature (and if, indeed, kinetic theory can be trusted anywhere for the QCD plasma). This, then, is the time needed to degrade the initial momentum of a charmed quark. It is quite a bit longer than the lifetime of the pure plasma phase. (If the initial temperature is as high as 500 MeV, the drag is much stronger and the lifetime is longer, so the effect is more significant.)

I have mentioned diffusion in momentum space, which in fact is the process of relaxation to equilibrium. The fluctuation-dissipation theorem gives the result

$$\langle (p - \langle p \rangle)^2 \rangle = m k_B T \left(1 - e^{-2\gamma t}\right) , \tag{11}$$

where γ is the same drag coefficient. Thus strong drag implies rapid thermalization.

If these estimates are to be believed, drag and thermalization for a charmed quark in the plasma is not a dramatic effect. The mixed phase, however, is quite another story. Here there is another source of drag, namely, *multiple hadronization.* Whatever the other characteristics of the plasma, quarks are not confined; confinement only sets in when they hadronize in the course of the phase transition. The phase transition, however, is a complex kinetic process, which proceeds through the nucleation, growth, and coalescence of bubbles of the low-temperature phase—the true QCD vacuum, populated with a gas of hadrons. This mixed phase is very long-lived, with perhaps ten times the lifetime of the original plasma. During this long time, a charmed quark may emerge again and again into bubbles of confining vacuum, only to reenter the plasma each time (see Fig. 1). Each emergence requires hadronization. In each hadronization, the charmed quark loses energy;[17] this *multiple* hadronization exerts a relatively large drag force on the quark, which will be reflected in the momentum of the charmed meson that is ultimately emitted. This drag thus offers a signal of the existence of a mixed phase in the course of the evolution of produced matter.

For a rough estimate of the drag force due to multiple hadronization, consider a charmed quark traveling through a region of mixed phase. On entering a bubble of vacuum, the charmed quark is forced to pull a light antiquark along to form a meson, losing energy in the process. (More precisely,[18] the charmed quark stretches a QCD flux tube into the vacuum, and the tube breaks either by sucking in an existing antiquark or by creation of a light $q\bar{q}$ pair.) Upon reentering the plasma, the light antiquark is stripped off and diffuses away, forcing the heavy quark to find another partner in the next hadronization. We estimate the rate of energy loss (the drag force) via

$$\frac{dE}{dx} = \frac{\langle \Delta E \rangle}{\lambda}, \tag{12}$$

where $\langle \Delta E \rangle$ is the mean energy loss in a single passage through a bubble and λ is the mean free path for quark-bubble collisions. The latter is given by

$$\lambda = (\rho\sigma)^{-1}, \tag{13}$$

where ρ is the density of bubbles and σ the cross-sectional area of a bubble, assumed the same for all bubbles at a given time.

A lower bound on $\langle \Delta E \rangle$ comes from the kinematics of the $c \to D \to c'$ hadronization/ionization process. Suppose the $c \to D$ step occurs at its kinematic limit, with no energy loss to the plasma or to the production of other particles. Then we have $E_D = E_c$, and the heavy quark is slowed because of the greater mass of the meson. Suppose similarly that the heavy quark loses no more momentum in the $D \to c'$ step, i.e., that it keeps its velocity as the light quark is stripped off. Then $E'_c/m_c = E_D/m_D$, and thus

$$\begin{aligned}\Delta E \equiv E'_c - E_c &= -\left(1 - \frac{m_c}{m_D}\right) E_c \\ &\approx -\frac{1}{6} E_c. \end{aligned} \tag{14}$$

The minimal value for this energy loss is obtained by setting $E'_c = m_c \approx 1.5$ GeV, namely, $|\Delta E| >$ 300 MeV. Thus even one extra hadronization causes substantial energy loss.

Creation of a D^* instead of a D, or fragmentation into $D\pi$, will more than double the energy loss. When the c quark is sufficiently energetic that threshholds are unimportant, these events are much more likely than D creation.

Returning to Eq. (13), we require estimates of ρ and σ. For these, we shall adopt a simple model of the phase transition kinetics. Begin with a plasma which cools to the transition temperature at proper time τ_p. (We work in the framework of the Bjorken scaling expansion.[3]) We suppose that at this instant there is enough fluctuation in the fluid to allow instant nucleation of bubbles of vacuum, with some density ρ_0; furthermore we assume that these bubbles can grow quickly enough to prevent any supercooling which would bring about additional nucleation at later times. Thus the number of bubbles is conserved, and the one-dimensional scaling expansion gives

$$\rho = \rho_0 \frac{\tau_p}{\tau}. \tag{15}$$

Since all the bubbles grow in parallel, they present identical geometrical cross-sections to the charmed quark, and we shall take $\sigma = \pi r^2$ to be the effective cross section in (13). $r(\tau)$ may be calculated from

$$\frac{4}{3}\pi r^3 \rho = 1 - R, \tag{16}$$

where $R(\tau)$ is the volume fraction of plasma in the mixed phase, which drops from unity (at τ_p) to zero (when the plasma disappears). R in turn can be calculated via entropy conservation, and the result is

$$1 - R = \frac{c_p}{c_p - c_h}\left(1 - \frac{\tau_p}{\tau}\right), \tag{17}$$

where c_p and c_h are the effective numbers of degrees of freedom in the respective phases. This gives us the values of $r(\tau)$ and $\rho(\tau)\sigma(\tau)$ plotted as solid curves in Figs. 2–3.

This relation is only valid, of course, as long as the bubbles don't touch. After the bubbles begin to coalesce, the surfaces join in such a way as to lower the total surface area. Coalescence *must* occur by the point at which $1 - R = 0.74$, which is the packing fraction for a face-centered-cubic lattice of dense-packed bubbles (an admittedly fanciful possibility).

After appreciable coalescence has taken place, we have an inverted picture: Now there are shrinking droplets of plasma surrounded by the hadron gas. The passage of the charmed quark through the medium at this stage is however similar to that discussed above, namely, a series of hadronizations alternating with ionizations. We continue using Eqs. (12)–(13), where now ρ represents the density of droplets and σ the droplet cross section. We suppose that the number of droplets is roughly equal to the original number of bubbles, so that (15) remains in force; replacing (16) with

$$\frac{4}{3}\pi r^3 \rho = R\,, \tag{18}$$

we obtain the dashed curves in Figs. 2–3, which go to zero at the end of the phase transition.

We see a period of peak drag lasting a time of perhaps $4\tau_p$ from the onset of the phase transition. Inserting the threshold value of $\Delta E = 300$ MeV, and guessing $\rho_0 = 1\ \mathrm{fm}^{-3}$, we obtain $|dE/dx| > 180$ MeV/fm during this period. Note that for an initial temperature of $T_0 = 1.5\,T_c$ at $\tau_0 = 1$ fm, the time scale in the figures is $\tau_p = 3.4$ fm, so the period of peak drag is long indeed.

Now let us consider the effect of this strong drag on the transverse momentum of the emerging charmed particles. We begin with pp collisions, where the distribution of D mesons is measured to be

$$\left(\frac{d\sigma}{dp_\perp^2}\right)_{pp\to DX} \propto e^{-p_\perp^2/p_{\perp 0}^2}, \tag{19}$$

with $p_{\perp 0} = 1.0$ GeV. A D meson traveling with a typical $p_\perp$ of 1 GeV has a velocity of $0.5\,c$ in the frame where its $\eta = 0$, and $E = 2.06$ GeV in that frame. We take this to be indicative of the energy of the parent c quark, and also of the typical c quark created in nucleus–nucleus collisions. Such a quark thus has $p_\perp = 1.4$ GeV and a kinetic energy of 560 MeV. This quark will stop after traversing only 3 fm! The collision region in a central Pb–Pb collision is some 15 fm in diameter, so it is clear that any quark created not too near the surface will stay within the collision region well into the mixed-phase stage, and furthermore will lose most of its initial momentum.

Can the c quark (or the D meson) be thermalized as it is stopped? If so, the transverse momentum distribution will be

$$\left(\frac{d\sigma}{dp_\perp^2}\right)_{AA\to DX} \propto e^{-p_\perp^2/2m_D T}. \tag{20}$$

The momentum scale appearing in (20), namely $\sqrt{2m_D T}$, will equal the scale $p_{\perp 0}$ appearing in (19) if $T = 280$ MeV as the D decouples from the medium. In that case the effect would disappear. If the D decouples at, say, $T = 200$ MeV, there would be a reduction of 20% in the momentum scale relative to that in the pp data.

But consider now the source of this thermalization. If it is interaction with the plasma, we have a calculation of the drag coefficient that shows it to be small—and it is this coefficient which governs the rate of thermalization according to (11). Thus thermalization by the plasma should be very weak. Furthermore, even if the plasma drag coefficient turns out to be much larger, any momentum imparted by the plasma will be lost to drag at the bubble surfaces in the mixed phase.

This brings us back to the bubbles. I don't know whether the bubble walls are in kinetic equilibrium with the matter they contain. (The *mathematical* applicability of the fluctuation-dissipation theorem is also questionable, since the drag force here is *not* proportional to the quark momentum as $\vec{p} \to 0$.) If the bubble walls are indeed cold, they will slow the charmed quarks without reheating them, giving very low $p_\perp$. If they are warm, the charmed quarks will end up thermalized at the transition temperature. In this case, charmed particles will constitute a precise *thermometer* of the transition temperature. The effect will still be dramatic if the transition temperature is below 200 MeV.

Finally, there is the pion gas which remains after the transition is over. I believe the drag due to pions at the transition temperature will be weak, since after all a pion gas is just a plasma with big holes in it (cf. the calculations of J/ψ suppression in the pion gas). Moreover, the pion gas will cool rapidly once the transition is over, so charm will decouple from the drag force quickly. The same argument applies to the question of whether a simple pion gas can mock up the effect by itself. Still, the influence of a dense hadronic medium should be estimated.

Two further questions come to mind. One, raised by Larry McLerran, is whether transverse hydrodynamic expansion can restore the charm $p_\perp$. I believe this to be unlikely, since both hydrodynamic calculations[20] and cascade simulations[21] show that transverse flow becomes significant only once the phase transition is well under way. The period of peak drag, on the other hand, occurs early in the phase transition, much of it before the bubbles coalesce into the rapidly diluted droplets.

The other question, raised by Gideon Alexander, is whether initial state scattering can enhance the charm $p_\perp$ in the same way as it does for the J/ψ. This could indeed be important. I can only hope that the J/ψ momentum can be used to gauge the effect of initial state interactions, which can then be subtracted off the momentum of the free charm. (As I once heard from Carlo Rubbia, "Yesterday's news is today's calibration, and tomorrow's background.")

SUMMARY

In order to motivate the introduction of a new signal of the quark-gluon plasma, I began by criticizing four existing signals. (I hope the parents of these ideas will forgive me for knocking their brainchildren.) Dileptons, strangeness enhancement, and J/ψ suppression all have practical problems, and in addition they depend on the notion of a weakly coupled, perturbative plasma for information about thermodynamics and correlations. Pion interferometry is a more robust tool, capable of observing the history and structure of the confinement phase transition without reference to details of the phases it separates; unfortunately, it, too, suffers from backgrounds and statistics.

Charm, in contrast, travels through the phase transition and gives a signal of that transition in its transverse momentum distribution. I believe this signal is *not* duplicated by any other dense, hot environment, but some calculations remain to be done. Finally, if it turns out that charm measurements are truly impossible, a new emphasis on observing the kinetics of the phase transition may still lead to a new family of signals.

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REFERENCES

1. B. Svetitsky, Tel Aviv preprint TAUP 1728-89 (unpublished, May, 1989). In the present talk I expand the discussion of rethermalization and its consequences.
2. The literature on photon and dilepton emission is vast. I have chosen to pick on the work of K. Kajantie, J. Kapusta, L. McLerran, and A. Mekjian, Phys. Rev. D **34** (1986) 2746. See references therein.
3. J. D. Bjorken, Phys. Rev. D **27** (1983) 140.
4. S. Hirasawa, M. Kadoya, and T. Miyazaki, Phys. Lett. **218B** (1989) 263.
5. T. Matsui and H. Satz, Phys. Lett. **178B** (1986) 416.
6. F. Karsch and R. Petronzio, Phys. Lett. **193B** (1987) 105.
7. See for example L. Kluberg, Nucl. Phys. **A488** (1988) 613c.
8. For reviews see J.-P. Blaizot, talk given at the 7th International Conference on Ultra-Relativistic Nucleus-Nucleus Collisions, Lenox, Massachusetts, September 1988, to appear in proceedings (Nucl. Phys. A, *in press*);
   S. Gavin, this volume.
9. S. Gavin, M. Gyulassy, and A. Jackson, Phys. Lett. **207B** (1988) 257;
   R. Vogt, M. Prakash, P. Koch, and T. H. Hansson, Phys. Lett. **207B** (1988) 263.
10. S. J. Brodsky and A. H. Mueller, Phys. Lett. **206B** (1988) 685.
11. A. Linde, Rep. Prog. Phys. **42** (1979) 389; Phys. Lett. **96B** (1980) 289;
    D. Gross, R. Pisarski, and L. Yaffe, Rev. Mod. Phys. **53** (1981) 43.
12. J. Rafelski and B. Müller, Phys. Rev. Lett. **48** (1982) 1066.
13. J. Kapusta and A. Mekjian, Phys. Rev. D **33** (1986) 1304;
    T. Matsui, B. Svetitsky, and L. McLerran, Phys. Rev. D **34** (1986) 2047.
14. G. Bertsch, M. Gong, and M. Tohyama, Phys. Rev C **37** (1988) 512;
    P. V. Ruuskanen and D. Seibert, Phys. Lett. **213B** (1988) 227.
15. But see A. Shor, Phys. Lett. **215B** (1988) 375.
16. B. Svetitsky, Phys. Rev. D **37** (1988) 2484.
17. For experimental data on charm fragmentation in what is perhaps an appropriate energy range, see P. Rapidis *et al.,* Phys. Lett. **84B** (1979) 507.
18. B. Banerjee, N. Glendenning, and T. Matsui, Phys. Lett. **127B** (1983) 453.
19. M. Aguilar-Benitez *et al.,* Phys. Lett. **189B** (1987) 476.
20. H. von Gersdorff, L. D. McLerran, M. Kataja, and P. V. Ruuskanen, Phys. Rev. D **34** (1986) 794;
    M. Kataja, L. D. McLerran, P. V. Ruuskanen,and H. von Gersdorff, Phys. Rev. D **34** (1986) 2755;
20. G. Bertsch, M. Gong, L. D. McLerran, V. Ruuskanen, and E. Sarkkinen, Phys. Rev. D **37** (1988) 1202.

Fig. 1. A charmed quark traversing a vacuum bubble as a D meson.

Fig. 2. Radii of vacuum bubbles (solid) and plasma droplets (dashed) in the mixed phase.

Fig. 3. Inverse mean free path for collisions with bubbles (solid) and droplets (dashed).