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SIGNATURES OF NEW PHENOMENA IN ULTRARELATIVISTIC NUCLEAR COLLISIONS

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Three classes of observables are discussed which may shed light on the properties of the quark-gluon plasma formed in ultrarelativistic nuclear collisons. They are (1) thermometers: the penetrating probes $\mu^+\mu^-$, γ, c, (2) barometers: transverse flow via $\langle p_\perp \rangle$, and (3) seismometers: fluctuations of dN/dy and $dE_\perp/dy$. The need for reliable estimates of the background due to the non-equilibrium processses is emphasized .

1. INTRODUCTON

In the past few years, much excitement has been generated by the realization that it is theoretically and technically feasible to create matter in the laboratory with energy density 10-100 times that found in ordinary nuclei[1,2]. Such energy densities are of interest because under such extreme conditions there may be a fundamental change in the properties of hadronic matter[3,4]. Current QCD estimates[5] indicate that strongly interacting, hadronic matter may dissolve into a weakly interacting quark-gluon plasma when the energy density exceeds $\varepsilon_Q \sim$ few GeV/fm^3. The possibility of producing and observing that new form of matter has also led to several proposals to construct a nuclear collider facility reaching center of mass energies up to 30 GeV/A on 30 GeV/A.

In the past year, the main theoretical progress has been the clarification of the expected initial conditions and space-time evolution of the plasma. On the experimental side, plans for detector systems capable of handling the expected 1000 charged particle multiplicities have been refined. The main challenge we face today is to connect the theoretical speculations with down to earth experimental signatures. That connection, as we shall see, is still tenuous. However, there are a few promising directions which warrant serious consideration. Of course, there remain many unressolved issues, and I aim to identify at least some of the key ones in the hope of stimulating more work in this area.

This talk is organized as follows: in Section 2, I review briefly the current understanding of the properties of quark-gluon plasmas. In Section 3, I emphasize the dynamical nature of the plasma produced in nuclear

collisions. Having indicated what we want to probe and in what environment we can expect to find it, I turn in Sections 4 thru 6 to specific observables and how they could serve as plasma diagnostic tools. In Section 3, the use of penetrating particles such as $\mu^+\mu^-$, γ, and charm as thermometers is discussed. Section 4 deals with barometers such as the average transverse momentum as a function of rapidity density. Finally, Section 6 describes seismometers, i.e., the study of fluctuations of, for example, $dE_\perp/dy$. These could provide signatures of explosive processes in the plasma. Concluding remarks are left for Section 7.

2. QCD NEWSREEL

The simplest estimate for the critical energy density at which a fundamental change in the properties of hadronic matter could occur comes from geometrical considerations. In normal nuclei, matter is rather dilute. There is only one nucleon per 7 fm^3, ($\rho_0 \cong .145\ fm^{-3}$), and the energy density is $\epsilon_{Nuc} = m_N\rho_0 \cong 0.15\ GeV/fm^3$. On the other hand, the energy density within a typical hadron is $\epsilon_H \sim 0.5\ GeV/fm^3 \sim 3\ \epsilon_{Nuc}$. Thus, with a modest increase of energy density from ϵ_{Nuc} to ϵ_H the dilute condition will change to a dense condition in which neighboring hadronic wavefunctions overlap and the internal degrees of freedom become activated. This increase in energy density can be achieved either by compressing cold nuclear matter or heating up the matter and filling the space between the nucleons with mesons. Thus starting in the hadronic world we can expect a marked change in the properties of matter when ϵ reaches $\sim\epsilon_H$.

From the asymptotic free QCD side,[3-5] we expect that at very high energy densities $\epsilon > \epsilon_Q$; the quark-gluon system behaves as a simple ideal (Stefan-Boltzmann) gas. Bag model estimates[5] give typically $\epsilon_Q \sim$ few $GeV/fm^3 \sim 10\ \epsilon_{Nuc}$ as the point beyond which the plasma phase is reached.

Such general considerations point to two characteristic numbers, ϵ_H and ϵ_Q, marking the transition region between our complex hadronic world and the simple quark-gluon plasma world. Figure 1 illustrates the qualitative dependence of ϵ and pressure, p, on temperature, T.

For more quantitative estimates, we must turn to the Monte Carlo lattice simulations of QCD. What is now well established[7,8] is that in pure gluon (Yang-Mills) QCD there is a strong first order phase transition beginning at $\sim\epsilon_H$ and ending at $\sim\epsilon_Q$, the critical temperature being $T_c \sim 200 \pm 50$ MeV. However, there is much uncertainty at present on the effect of including quarks into such calculations. Recently DeGrand and DeTar[9] showed that in a model Z(3) theory the inclusion of light fermion loops and/or high chemical

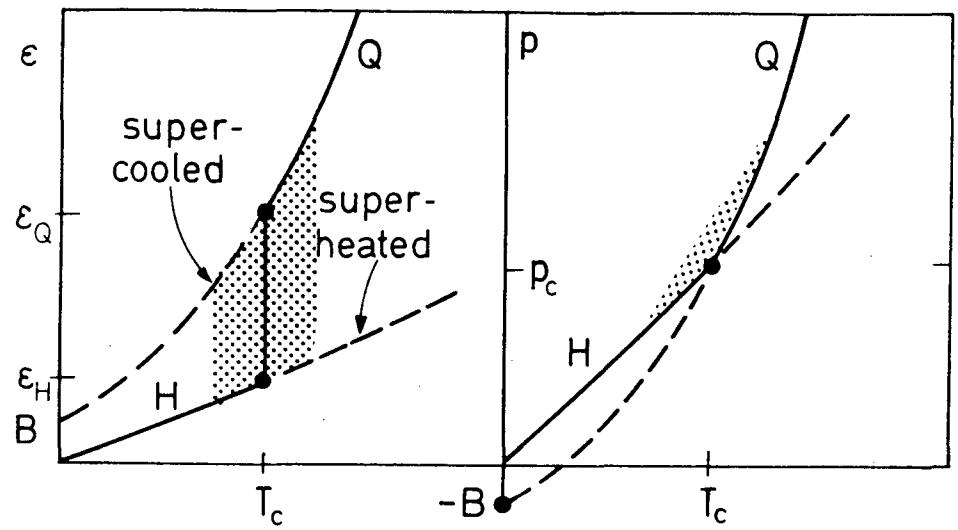

FIGURE 1
Energy density and pressure as a function of temperature. In the shaded region around $T_c \sim 200 \pm 50$ MeV there is a rapid change of the temperature dependence as hadrons dissolve into a quarks-glue plasma and chiral symmetry is restored. Solid curves show Bag model equation of state.[5]

potential can wash out the first and even second order character of the phase transition. On the other hand, the Bielefeld group[10] showed with another approximate treatment of fermion loops that while the singularities of $\varepsilon(T)$ are washed out, the specific heat, $\partial\varepsilon/\partial T$, is still peaked between ε_H and ε_Q. Additional confusion (entropy) has been generated by a recent report[11] that the lattice Λ parameter may be twice as large as previously assumed. These issues remain "hot" topics and reflect the theoretical uncertainties in the nature of the hadron to plasma transition. From the point of view of observables, it is obvious that we should only consider generic observable that have flexibility to accommodate a variety of possible phenomena.

There is yet another level of uncertainty we must keep in mind. That concerns the very nature of the plasma[12]. Assuming that after the dust has settled the value of ε_Q remains on the order of 10 - 100 ε_{Nuc}, what can be assumed about the properties of that plasma? The value of ε_Q would be determined by where the energy density becomes well approximated by the ideal Stefan-Boltzmann law, $\varepsilon = KT^4$, where $K \cong 12.2$ is the constant determined by the number of helicity, flavor, and color states in the plasma. The claim of Carruthers[12] is that a lattice "measurement" of KT^4 behavior is not enough to guarantee that the plasma is ideal. His point is that the quasiparticles in the plasma may be phonons and plasmons rather than quarks and gluons. The

plasmons would arise because near the transition temperature the effective couplings α_s could be large and a dynamical mass $m_g \sim T_c$ could be generated by gluons. At first sight such a possibility would appear to be ruled out by determination of $K \cong 12.2$ as expected if $m_g = 0$.

If there were a gluon mass, then the energy density of gluons would be reduced from $\epsilon_g(0) = 8\pi^2T^4/15$ to $\epsilon_g(m)$, where for $3m/2 > T > m/3$ an approximate form is given by[13]

$$\frac{\epsilon_g(m)}{\epsilon_g(0)} \cong \frac{1}{3}\left(\frac{T}{m}\right)^{7/2}\left(\frac{200\ \text{MeV}}{T}\right)^3 . \qquad (2.1)$$

Thus for $m_g \sim T_c \sim 200$ MeV, the Boltzmann factor could reduce the gluon contribution to the energy density by a factor ~1/3. Note the sensitivity of eq. (2.1) to $m_g(T)$. Hence, K should be smaller or larger than expected. However, additional energy density may be tied up in excitations of ordinary <u>first</u> sound modes in the plasma. A simple estimate of the contribution from such phonons can be made as follows: Sound waves can remain undamped down to wavelengths, $\lambda = 2\pi/k$, on the order of the mean free path λ_{mfp} in the system. Near T_c, λ_{mfp} is probably as small as the uncertainty principle allows[14], i.e., $\lambda_{mfp} \sim \hbar/T_c$. Thus, sound modes with wavevectors up to $k < 2\pi T_c$ could propagate undamped. Their contribution to the energy density would be

$$\epsilon_{ph} \cong \int_{k<2\pi/\lambda_{mfp}} \frac{d^3k}{(2\pi)^3}\,\frac{\omega}{e^{\omega/T}-1} \cong \frac{\pi^2}{30}\,\frac{1}{c_o^3}\,T^4 , \qquad (2.2)$$

where we used dispersion relation for sound, $\omega = c_o k$, $c_o^2 = 1/3$, and took $k_{max} = \infty$ on grounds that the integrand peaks below $k \sim 2\pi T$. Carruthers thus found for $T \sim T_c$,

$$\epsilon_{ph} \sim \frac{1}{3}\,\epsilon_g\,(m_g = 0) \qquad (2.3)$$

Therefore, much of the lost energy density in eq. (2.1) could be compensated for by first sound. This could lead to "precocious" Stefan-Boltzmann behavior of ϵ!

It is amusing to note that for ordinary ideal gases $2\pi/\lambda_{mfp} << T$, and thus $\epsilon_{ph} \cong (4/3)\pi T/\lambda^3$. The ideal gas contribution $\epsilon = 3/2\ T\rho$ is in this case larger by a factor $\epsilon/\epsilon_{ph} \sim \rho\lambda_{mfp}^3 \sim 10^6$ at STP. Therefore, unlike a quark-gluon plasma, ordinary ideal gases are not very noisy. The plasma may be more similar to hydrogen gas at pressures $\sim 10^3$ atm.

The point of this exercise is to alert us to the possibility that $\epsilon \cong KT^4$ is not enough to show that the relevant quasiparticles are free quark and gluon states. Again, from the point of view of observables and signatures such possibilities must be kept in mind. In particular, model calculations assuming that the plasma is ideal could lead to erroneous expectations.

3. THE DYNAMIC PLASMA

The challenge of finding diagnostic tools to study the quark-gluon plasma is heightened not only by the uncertain nature of the plasma and its transition back to the hadronic world, but also by the dynamic environment in which it is produced.

3.1. The scaling regime, $E_{lab} > 1$ TeV/A

Consider first the low baryon density plasma that we expect to produce in the central region at energies $E_{cm} > 30$ GeV/A. At sufficiently high energies, the fragmentation regions, containing the baryons, separate leaving a region of rapidity space occupied mainly by mesons. That region is characterized by a approximately constant rapidity density, dN/dy. As emphasized by Bjorken,[15] the constancy of dN/dy means that the evolution of the plasma is invariant under Lorentz boosts along the beam axis. Physically this means that field variables such as energy densities, ϵ, pressures, p, and entropies, s, can only depend on proper time, $\tau = (t^2 - z^2)^{1/2}$. Thus, in a space-time diagram, contours of constant ϵ, p, σ correspond to simple hyperbolic lines τ = const. In addition, the flow velocity of the plasma at any point (t,z) can be computed as

$$v_z(t,z) = \tanh y = z/t \quad . \tag{3.1}$$

From eq. (3.1), we see that the "Hubble" constant for the plasma

$$H = \frac{\partial v_z}{\partial z} = \frac{1}{t} \sim \left(\frac{fm}{c}\right)^{-1} \tag{3.2}$$

is large. In fact, this Hubble constant is about 10^{17} larger than the cosmological Hubble constant at the time when the temperature of the universe was $T \sim 200$ MeV. Thus, while the low baryon density plasma produced in

nuclear collisions is similar to the primordial Big Bang plasma at times $t \sim 10^{-6}$ sec, there is this very important dynamical difference. The plasma in nuclear collisions is created with a tremendous longitudinal velocity gradient, eq. (3.2). This gradient is in fact as large as it could possibly be according to the uncertainty principle and relativity. The value of the initial time $\tau_0 \sim 1$ fm/c comes frm the uncertainty principle applied to the emission of finite $p_\perp \sim m_\pi$ quanta. For $\tau < \tau_0$ the plasma is intrinsically in a quantal state. Before $\tau \sim \tau_0$ it makes no sense to talk about temperatures $T \simeq 200$ MeV because energies uncertainties $\hbar/\tau$ exceed T. Only after $\tau > \tau_0$ can we begin to speak of classical thermodynamic and hydrodynamic phenomena. In comparison, the analogous quantal time, when $H \sim T$, for the Big Bang is $\tau_0 \sim 10^{-42}$ sec.

This large longitudinal velocity gradient has important consequences. Most important is the rapid cooling that results[15,16]. That rate of cooling is in fact <u>independent</u> of the nuclear dimensions! To see this recall Bjorken's analysis. First, consider the case where there are no final state interactions and the plasma expands freely. Because of eq. (3.1) the volume element of the plasma containing quanta of rapidities $-\partial y < y < \partial y$ increases linearly with time. Since energy is conserved, the energy density must then decrease as

$$\varepsilon(\tau) = \varepsilon_0 \left(\frac{\tau_0}{\tau}\right) \quad : dE = 0 \quad . \tag{3.3}$$

The symbol $dE = 0$ is to remind us that eq. (3.3) is derived under the assumption of iso-ergic expansion. In contrast, if sufficiently strong final state interactions occur to generate local equilibrium, then the expansion proceeds isentropically ($dS = 0$). In that case, $\partial_\mu T^{\mu\nu} = 0$ reduces to

$$\frac{d\varepsilon}{d\tau} + \frac{\varepsilon + p}{\tau} = 0 \quad . \tag{3.4}$$

Therefore, for an equation of state $p = c_0^2 \varepsilon$

$$\varepsilon(\tau) = \varepsilon_0 \left(\frac{\tau_0}{\tau}\right)^{1+c_0^2} \quad : dS = 0 \quad . \tag{3.5}$$

The reason that ε falls faster in $dS = 0$ expansion is that pdV work is being done upon expansion[17]. That energy propagates ultimately into the fragmentation regions.

The main lesson to learn from eqs. (3.3-3.5) is that within the first few fm/c, the energy density will fall by a large factor. To appreciate this rate of expansion note from Fig. 1 that if we start at $\tau_0 \sim 1$ fm/c with $\varepsilon > \varepsilon_Q$, then by $\tau \sim 2$ fm/c, the energy density has decreased to a point well into the transition region between the quark and hadron worlds. However, in order for the plasma to evolve adiabatically, the reactions rates Γ must exceed the expansion rates H. On the other hand, Danielewicz[14] recently showed that in systems near equilibrium the rates are bounded by $T/\hbar$ due to the uncertainty principle. This result can be understood as follows: for a given Γ the uncertainty in single particle energies must exceed $\Delta E > \Gamma$. To be near equilibrium requires that $\Delta E < T$. Putting these results together a necessary condition for adiabatic evolution of the plasma is

$$T/\hbar > \Gamma \gg H \quad . \tag{3.6}$$

Assuming that is as large as it can be, that H is given by eq. (3.2), and that $T(\tau) = T_0(\tau_0/\tau)^{1/3}$ according to eq. (3.5), eq. (3.6) reduces to

$$\frac{T_0 \tau_0}{\hbar} \left(\frac{\tau}{\tau_0}\right)^{2/3} \gg 1 \quad . \tag{3.7}$$

Again by the uncertainty principle, the earliest time τ_0 at which it is sensible to talk about thermal equilibrium is $\tau_0 \sim \hbar/T_0 \sim 1$ fm/c. What eq. (3.7) shows is that after several fm/c, adiabatic evolution is possible. However, in the first few fm/c non-equilibrium phenomena are likely. In particular, it is possible that the plasma will find itself in a strongly supercooled[18] rather than mixed quark hadron phase.

We have emphasized that the dominance of longitudinal expansion is independent of nuclear size. However, actual signatures could depend on transverse surface radiation[19,20] and rarefaction[21] phenomena as well. For small systems, such as pp, where the transverse dimension, R, is comparable to the mean free paths, $\lambda_{mfp} > \hbar/T$, non-equilibrium transverse expansion could prevent the formation of any locally equilibrated phase. In the largest systems, UU, the transverse rarefaction wave[21] can lead to more rapid cooling of the outer portion of the plasma, and only the interior core of the plasma would evolve according to scaling dynamics. The finite transverse size also implies that the initial energy densities, $\varepsilon(\tau_0)$, cannot be uniform with respect to the transverse coordinate,

$\varepsilon \propto (R^2 - x_\perp^2)^{1/2}$. Finally, for finite energies the local baryon density in the central region will depend on transverse coordinate. In general, regions of highest energy density are associated with highest baryon density because both quantities grow with increasing nuclear thickness.

These considerations obviously imply that all observables involve a complex convolution over the space-time history of the reaction. There is no simple spherical fireball produced. The only hope then of unfolding all those convolutions will be to cross check as many different observables as possible and study carefully systematics with respect to variations in atomic number, energy, multiplicity.

3.2. The stopping regime, $E_{lab} \sim 10$ GeV/A

We discuss next the environment of the high baryon density plasma[22-25] produced at lower energies. Longitudinal growth, $z = \gamma\tau_0$, which follows from the uncertainty principle and special relativity, leads to formation distances z that exceed the nuclear dimensions when lab energies exceed,[25] $E_{lab} > 2R/\tau_0$ GeV/A ~ 10 GeV/A for uranium. This leads to nuclear transparency and eventually to scaling dynamics. However, below 10 GeV/A, two uranium nuclei can stop each other in the center of mass system. The stopped matter can reach energy densities $\varepsilon_0 \sim$ few GeV/fm^3 as in the scaling regime, but the baryon densities can be as high as 10 ρ_0. Figure 2 constrasts the space-time evolution of such reactions with those in the scaling regime. In the cm, two Lorentz contracted nuclei with thickness, $2R/\gamma_{cm}$, begin to interact at t = 0. The fraction of initial nucleons that can contribute to direct or knockout reactions is given by $f = \tau_0\gamma_{cm}^2/2R$, < 1 in the stopping regime.

The direct reactions occur in the relatively small space-time volume shaded in Fig. 2. This is in contrast to the scaling regime, where in the first fm/c, <u>all</u> partons of both nuclei pass through one another and have a chance at a direct interaction. This difference implies that direct processes such as Drell-Yan could be a much larger source of background in the scaling regime than in the stopping regime.

The second major difference is that the intermediate region, marking the space-time domain where increasing number of collisions are driving the system toward local equilibrium, is less important for E < 10 GeV/A. The thickness of that region is the longitudinal growth length,[25] $z = \gamma_{cm}\tau_0$. It extends in time until $t \sim R_{cm}/\gamma_{cm}$. In that region, the incoming kinetic energy is being converted into heat and compression, and thus represents a shock front. The highest energy densities are reached in this case in the

thermal region between the expanding shock fronts. In contrast, in the scaling regime the highest energy densities occur in the intermediate region.

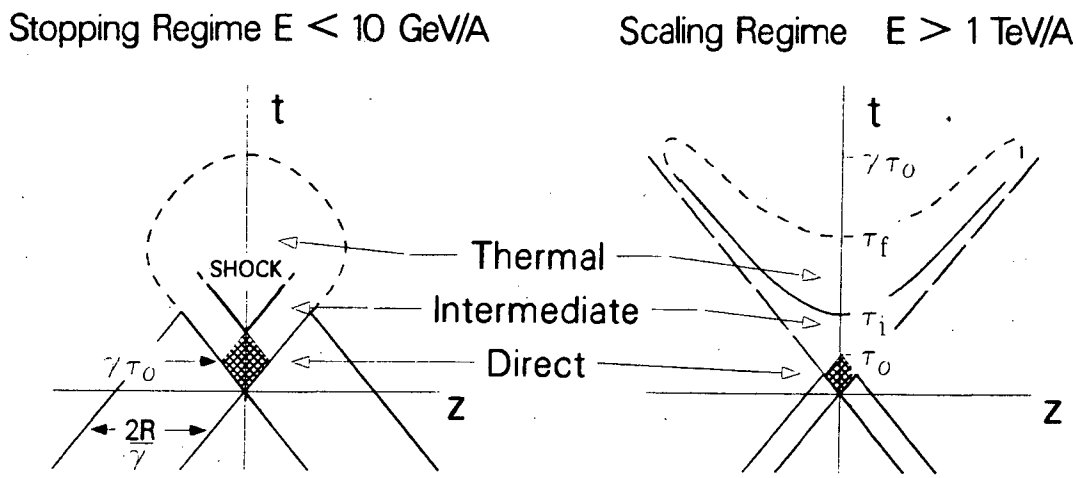

FIGURE 2
Space-time evolution of the quark-glue plasma in the stopping and scaling regimes. Two Lorentz contracted nuclei collide in the center of mass with an energy γm_N per nucleon. Three regions corresponding to direct reactions, the approach toward local equilibrium, and hydrodynamic expansion are contrasted in the two regimes.

Finally, another major difference is that the shocked matter is relatively stationary with respect to temperatures and densities until $t \sim R/\gamma_{cm}$. Thus the period of highest energy densities is prolonged in the stopping regime relative that in the scaling regime.

In summary, the dynamical history of the high baryon density plasma produced at 10 GeV/A is very different from that of the low baryon density plasma produced above 1 TeV/A. Obviously, interpretations of signatures will have to take this difference into account.

4. THERMOMETERS

Given the preceding remarks, we proceed to consider specific observables. One of the first things we would want to know is whether local thermal equilibrium had been established and what the thermal history was. As with deducing the interior temperature of the sun, penetrating probes would probably be lest in this regard. For quark-gluon plasmas, dilepton pairs $\mu^+\mu^-$, direct photons, γ, or heavy flavors, c, have been suggested as approriate thermometers[26-29]. I will follow Ref. (26) and concentrate on

how the dynamic nature of the plasma affects the calibration of such thermometers.

Let $\Gamma(x)$ be the rate per unit volume to produce one of those penetrating probes. That rate depends on the local distribution of quarks and gluons. For example, if $f_q(x,p)$ and $f_{\bar{q}}(x,p)$ are the (Wigner) densities of quarks and antiquarks at space-time x and three momentum p and p' then the dilepton rate has the form

$$\Gamma(x) = \int d^3p\, d^3p' \, [v\sigma_{q\bar{q}\to l\bar{l}}]\, f_q(x,p)\, f_{\bar{q}}(x,p'), \tag{4.1}$$

where $v\sigma$ is the annihilation rate. For x in the thermal regions in Fig. 2, f_q and $f_{\bar{q}}$ reduce to simple Fermi-Dirac functions, $(e^{\omega/T}+1)^{-1}$. In that case $\Gamma(x) = \Gamma(T(x))$, where $T(x)$ is the local temperature.

However, at early times f_q, $f_{\bar{q}}$ are very far from equilibrium. In fact, in the shaded region in Fig. 2, they measure only the initial distributions of quarks and antiquarks and are related to nuclear structure functions. If it were not for final state interactions, the penetrating probes would be exclusively produced from the direct reactions occuring during that initial time. This contribution corresponds for dilepton pairs to the usual Drell-Yan yield.

In addition to the direct and thermal regions, there is an important non-equilibrium intermediate region in space-time as shown in Fig. 2. In that region, secondary quanta are being produced and collisions among and with them are leading to local equilibration. This is obviously the most complex stage of the reactions and requires a transport theory to specify the evolution of f_q and $f_{\bar{q}}$ from the initial structure functions toward the Fermi-Dirac form.

Taking into account all three stages, the yield of penetrating particles can be decomposed as

$$N = \int d^4x\,\Gamma(x) = N_{Dir} + N_{Int} + N_{Th}$$

$$= \int_{\Omega_{Dir}} d^4x\ \Gamma(x) + \int_{\Omega_{Int}} d^4x\ \Gamma(x) + \int_{\Omega_{Th}} d^4x\ \Gamma(T(x)), \tag{4.2}$$

where Ω_{Dir}, Ω_{Int}, and Ω_{Th} are the four volumes in Fig. 2 corresponding to direct, intermediate and thermal stages of the reaction. One should also add the yield associated with decays of final hadrons (ω,η,π^0). It is the

last term in eq. 4.2 that we are interested in. Defining the thermal profile distribution[26]

$$\Phi(T) = \int_{\Omega_{Th}} d^4x \; \delta(T - T(x)), \tag{4.3}$$

the thermal contribution is

$$N_{Th} = \int dT \; \Gamma(T) \; \Phi(T). \tag{4.4}$$

By studying differential yields (per invariant mass or $p_\perp$) and looking at different probes, $\mu^+\mu^-$, γ, , the end objective would be to deduce $\Phi(T)$.

Before discussing the direct and intermediate background, consider the expected form of $\Phi(T)$. In the scaling regime $\Omega_{Th} = 1$ for $\tau_i < \tau < \tau_f$, where $T(\tau_i) = T_o$ and $T(\tau_f) = T_f$, the "freezeout" temperature. In terms of the rapidity variable, y, defined in eq. (3.1), the thermal region is bounded by $-Y_{cm} < y < Y_{cm}$, with Y_{cm} being the incident rapidity of nuclei in the cm system. Finally, transverse coordinates are restricted to $x < R$. For an ideal expansion, $p = \epsilon/3$, $T(x) = T(\tau) = T_o \; (\tau_i/\tau)^{1/3}$ in the scaling regime[15,16]. Therefore, the integral in eq. (4.3) gives with $d^4x = \tau d\tau dy d^2x_\perp$

$$\Phi(T) = \pi R^2 \; 2Y_{cm} \; \tau_i^2 \frac{1}{T_o} \left(\frac{T_o}{T}\right)^7 \theta(T_f < T < T_o). \tag{4.5}$$

Note that $\Phi(T)$ is heavily weighted toward <u>lower</u> temperatures because the rapid longitudinal expansion allows the system to remain at high temperatures only for short times.

The most important point to note in eq. (4.5) is that $\Phi \propto R^2 \propto A^{2/3}$. The effective four volume is not R^4 as initially guessed[27-29], but rather only $R^2\tau_i^2$. As emphasized by Kajantie, et al.[16] this consequence of longitudinal expansion can severely reduce the thermal signal from the direct and intermediate noise. This follows because N_{Dir} is the number of nucleon pairs interacting only once times the probability, $\sigma_{q\bar{q}\to l\bar{l}}/\sigma_{reac}$, that a penetrating probe is made during that direct reaction. Since $\sigma_{reac} \propto R^2$,

$N_{Dir} \propto A^2/R^2 \propto A^{4/3}$. Therefore, in the scaling regime, the direct component can overwhelm the thermal one for heavy nuclei.

In contrast, in the stopping regime the fraction of pairs interacting for the first time is only a small subset of total as seen in Fig. 2. The total number of possible direct NN reactions is proportional to A^2f^2 where $f = \tau_0\gamma_{cm}^2/2R$. Therefore, $N_{Dir} \propto A^2/R^4 \propto A^{2/3}$ in this regime. As a crude estimate, for $\Phi(T)$, note that the temperature and density are approximately constant in the shock wedge, $|x| < t - \gamma_{cm}\tau_0$ and $\gamma_{cm}\tau_0 < t < R/\gamma_{cm}$. Therefore,

$$\Phi(T) \cong \pi R^2\left(\frac{R}{\gamma_{cm}} - \gamma_{cm}\tau_0\right)^2 \delta(T - T_0), \tag{4.6}$$

is peaked at the temperature T_0 of the shocked matter. Also below the transparency limit $\gamma_{cm} \lesssim (R/\tau_0)^{1/2}$, $\Phi \propto R^4 \propto A^{4/3}$. Based on this, it appears that penetrating probes are better thermometers of the baryon rich plasma produced in the stopping regime than of the baryon free plasma produced in the scaling regime. However, we must remember that f_q and $f_{\bar{q}}$ depend also on the unknown chemical potential, $\mu(x)$, in the stopping regime. It would take a combination of measurements to determine T and μ separately. Such a measurement would of course be equivalent to measuring the equation of state of high density $\rho \sim 5\text{-}10\ \rho_0$, high temperature $T \sim 100\text{-}200$ MeV nuclear matter. The difficulty of such a measurement is well appreciated already from the study of 1 GeV/A nuclear collisions[30].

In both regimes, the ultimate accuracy of the thermometer rests on our ability to estimate the intermediate component. This will require the further development of transport theories and Monte Carlo parton cascade simultations[31]. However, a few general remarks can be made at this time. Consider dilepton yields as a function of invariant mass as a specific example (see p. 475, 476 of Ref. 1). For very massive, $M > 4$ GeV, pairs, $N_{Th} \sim e^{-M/T}$ dies very fast in comparison to the power law behavior, $N_{Dir} \sim M^{-4}$, expected for Drell-Yan. For low mass pairs, $M < 400$ MeV, Dalitz pairs can dominate the yield. Shuryak suggested[33] that intermediate mass $m_\rho < M < 2$ GeV pairs would be ideal. However, in this region the non-equilibrium contribution must surely be important. Shuryak in fact needed a very large initial temperature $T \sim 500$ MeV to fit the intermediate mass dilepton yield. Such high temperatures (arising from the Landau boundary condition) are however incompatible with longitudinal growth unless the τ_0 parameter is less than 0.1 fm/c. Furthermore the data refer to hadron-nucleus collisions

where, due to finite $x_\perp < 1$ fm/c of the reaction zone, thermalization is much less likely. The intermediate mass pairs can be easily produced, on the other hand, during the non-equilibrium evolution of the plasma as the large relative longitudinal momentum of partons are being degraded by multiple collisions. It would be clearly desirable to perform estimates along such lines using a generalization of the cascade model developed in Ref. (31).

Finally, we want to comment on a recent calculation[32] on the effect of non-trivial dispersion relations, $\omega(k)$, of quarks in non-ideal plasmas. As noted in Section 2, it is possible[12] that the quasiparticles of the plasma produced in nuclear collisions are massive quarks and gluons with $m_g \sim m_q \sim T$. Instead of the ideal dispersion, $\omega = k$, it could be that $\omega = (k^2 + T^2)^{1/2}$ or a more complex form[32]. It is clear that the most sensitivity to such dispersion relations would be for the pair mass range $M < 2T < 400$ MeV. Kapusta showed that the mass spectrum of e^+e^- pairs of zero total momentum can vary by orders of magnitude in this low mass range. Unfortunately, Dalitz pairs from $\pi^\circ \to \gamma e^+e^-$ also expected[32] to dominate the thermal yield in that range by a factor $\sim 10^3$. In addition, because of the finite lifetime of the plasma and the rapid expansion rate, the dispersion relation is not well defined by the uncertainty principle for $\omega < \hbar H$, where H is given by eq. (3.2). During the hotest phase, $H \sim T_0/2$ in the scaling regime. This limits the range of sensitivity to $\omega > 100$ MeV, i.e. for $T_0 < M_{e^+e^-} < 2T_0$. In this regard, the stopping regime offers the advantage of providing a larger duration where temperatures and densities remain approximately constant. The main lesson to learn from this calculation is that in the low mass or $p_\perp$ region, M or $p_\perp < 2T_0$, the penetrating probes are not reliable thermometers. They could, on the other hand, be exploited to give useful information on the quasiparticle degrees of freedom in the plasma.

In conclusion, we see that the value of penetrating probes as thermometers depends sensitively on the kinematic domain in which they are measured. It appears that the range $1 < (M, P_\perp) < 2$ GeV is probably the best. However, the absolute calibration of these thermometers depends critically on our eventual ability to calculate the yields arising from the non-equilibrium stage of the reaction.

5. BAROMETERS

In addition to the temperatures, the pressures generated in the plasma are of key interest. Here we consider a recent suggestion[33,34] that the average transverse momentum, $\langle p_\perp \rangle$, may provide such a barometer. To see this

connection we recall[35] the theory of one dimensional simple waves. A simple wave is one where all components of the energy-momentum tensor, $T_{\mu\nu}$, can be expressed as functions of one variable, say the local energy density, ε. In that case the equations of motion, $\partial_\mu T^{\mu\nu} = 0$, reduce to

$$\frac{\partial\varepsilon}{\partial t}\frac{dT_{00}}{d\varepsilon} + \frac{\partial\varepsilon}{\partial x}\frac{dT_{01}}{d\varepsilon} = 0 \quad ,$$

$$\frac{\partial\varepsilon}{\partial t}\frac{dT_{01}}{d\varepsilon} + \frac{\partial\varepsilon}{\partial x}\frac{dT_{11}}{d\varepsilon} = 0 \quad , \tag{5.1}$$

from which it follows that[35]

$$\left(\frac{dT_{01}}{d\varepsilon}\right)^2 - \left(\frac{dT_{00}}{d\varepsilon}\right)\left(\frac{dT_{11}}{d\varepsilon}\right) = 0 \tag{5.2}$$

Utilizing the form of $T^{\mu\nu}$ for a perfect fluid

$$T^{\mu\nu} = (\varepsilon + p)\, u^\mu u^\nu - p g_{\mu\nu} \quad , \tag{5.3}$$

where the fluid velocity $u^\mu = (\cosh\eta, \sinh\eta)$ in terms of the fluid rapidity η, eq. (5.2) reduces to the simple relation

$$\frac{d\eta}{d\varepsilon} = -\frac{c_0}{\varepsilon + p} \quad , \tag{5.4}$$

where $c_0^2 = dp/d\varepsilon$ is the sound velocity squared. Given the equation of state $p(\varepsilon)$, we can integrate eq. (5.4) to yield the fluid rapidity when the energy density is reduced from ε_0 to ε as

$$\eta(\varepsilon) = \int_{\varepsilon}^{\varepsilon_0} \frac{c_0(\varepsilon)d\varepsilon}{\varepsilon + p(\varepsilon)} \quad . \tag{5.5}$$

For an ideal gas, $p = c_0^2\varepsilon$, and therefore

$$\eta(\varepsilon) = \frac{c_0}{1 + c_0^2}\ln\left(\frac{\varepsilon_0}{\varepsilon}\right) = \frac{c_0}{1 + c_0^2}\ln\left(\frac{p_0}{p}\right) \quad . \tag{5.6}$$

Applying, eq. (5.6) to transverse expansion (neglecting curvature) shows that a measurement of the transverse flow velocity at the brake-up time (when $\varepsilon = \varepsilon_f \cong (0.1 - 0.5)$ GeV/fm^3 and $p = p_f$) could serve as a barometer of the initial pressure, p_0.

The simplest observable that depends on transverse flow is the average transverse momentum, $\langle p_\perp \rangle$. However, $\langle p_\perp \rangle$ depends not only on η but also on the internal temperatures, $T \sim m_\pi$, of the fluid at brake-up:

$$\langle p_\perp \rangle = m_\pi \sinh\left(y_T + \eta(\varepsilon_f)\right) \quad , \tag{5.7}$$

where $y_T \cong 1.5$ for $T \cong m_\pi$ is the average transverse rapidity due to random thermal motion. Because $y_T > 1$, we can approximate eq. (5.7) as

$$\langle p_\perp \rangle = p_T e^{\eta(\varepsilon_f)} = p_T \left(\frac{\varepsilon_0}{\varepsilon_f}\right)^{\frac{c_0}{1 + c_0^2}} = p_T \left(\frac{p_0}{p_f}\right)^{\frac{c_0}{1 + c_0^2}} , \tag{5.8}$$

where $p_T = m_\pi e^{y_T/2}$ 0.3 GeV/c. Eq. (5.8) shows that the calibration of the $\langle p_\perp \rangle$ barometer depends on knowledge of the speed of sound, c_0, and the pressure, p_f, at brake up.

A consistency check of the above relation is given by the dependence of $\langle p_\perp \rangle$ on the rapidity density. In the scaling regime an estimate of the initial energy density as a function of the final pion rapidity density is given by[17]

$$\frac{\varepsilon_0}{\varepsilon_f} \cong \left(\frac{3m_\pi}{\varepsilon_f \tau_0 A_\perp} \quad \frac{dN_\pi}{dy}\right)^{1 + c_0^2} , \tag{5.9}$$

where $A_\perp$ is the transverse area of the plasma at the time $\tau_0 \sim 1\text{-}2$ fm/c marking the onset of hydrodynamic expansion. Consequently, we get

$$\langle p_\perp \rangle \propto \left(\frac{dN_\pi}{dy}\right)^{c_0} , \tag{5.10}$$

Looking at pp data[36] this predicts that $\langle p_\perp \rangle = 0.8$ GeV/c for $dN_\pi^{ch}/dy = 10$ given that $\langle p_\perp \rangle = 0.36$ for $dN_\pi^{ch}/dy = 2.5$, in clear contradiction with the data where $\langle p_\perp \rangle < 0.5$ GeV/c. A smaller sound velocity, $c_0^2 = 1/6$, will not help either. What's wrong with the barometer?

Up to now we have neglected the effects of longitudinal expansion. As shown by Baym et al.[21] inclusion of this dynamical aspect of the plasma reduces greatly the transverse flow. To gain insight into how this comes about, note that the expansion proceeds via a rarefaction wave traveling inward at the speed of sound. The fluid at radius r will therefore not feel transverse pressure gradients before $t = \tau_0 + (R - r)/c_0$. But by that time the energy density of that fluid element will have decreased because of logitudinal expansion to

$$\epsilon_0(r) = \epsilon_0 \left(\frac{\tau_0}{\tau_0 + (R - r)/c_0}\right)^{1 + c_0^2} . \tag{5.11}$$

Therefore, smaller pressure gradients are available to accelerate that fluid element in the transverse direction.

For a rough estimate of the transverse rapidity acquired by the fluid element initially at radius r, we simple modify eq. (5.5. 5.6, 5.8) by replacing ϵ_0 by $\epsilon_0(r)$ given by eq. (5.11). In that way the average transverse momentum of particles arising from the fluid element initially at r is estimated to be

$$p_\perp(r) \cong p_T\left(\frac{\epsilon_0(r)}{\epsilon_T}\right)^{\frac{c_0}{1 + c_0^2}} \cong \langle p_\perp \rangle_0 \left(\frac{1}{1 + (R - r)/c_0\tau_0}\right)^{c_0} , \tag{5.12}$$

where $\langle p_\perp \rangle_0$ is given by eq. (5.8). Note that eq. (5.12) applies only for $r > r_c$ where $\epsilon_0(r_c) \equiv \epsilon_f$:

$$r_c = R - c_0\tau_0\left\{\left(\frac{\varepsilon_0}{\varepsilon_f}\right)^{\frac{1}{1+c_0^2}} - 1\right\} , \qquad (5.13)$$

as given in Ref. (21). For $r < r_c$, $p_\perp(r) = p_T$ because the fluid breaks up before the ratefaction wave can hit it.

Thus only a fraction of the plasma can acquire transverse flow in the first place. For sufficiently large ε_0, $r_c < 0$, but even then only the outer portion of the plasma acquires significant transverse flow.

Averaging over the initial transverse coordinate, the average transverse momentum is

$$\langle p_\perp\rangle = \int_0^R p_\perp(r)\,\frac{2rdr}{R^2} = p_T\left(\frac{\varepsilon_0}{\varepsilon_f}\right)^{\frac{c_0}{1+c_0^2}} F\left(\frac{\varepsilon_0}{\varepsilon_f}, \frac{R}{c_0\tau_0}\right), \qquad (5.14)$$

where F is a finite size form factor including effects of longitudinal expansion such that $F \to 1$ for $R \to 0$. Asymptotically as $R \to \infty$, $F \to (\varepsilon_f/\varepsilon_0)^{c_0/1+c_0^2}$, i.e. $\langle p_\perp\rangle \to p_T$. While the integral is analytic we prefer to show the results in Fig. 3. For R = 7 fm and $\varepsilon_0/\varepsilon = 4$, $T_f/T_0 \cong .7$), there is about a factor 10 reduction of the average transverse rapidity $\langle\eta\rangle$ from 0.6 for F = 1 to 0.07. Even for $\varepsilon_0/\varepsilon_f \cong 40$ ($T_f/T_0 \cong 0.4$) $\langle\eta\rangle$ is reduced from 1.6 to 0.77. These results are in qualitative accord with the detailed numerical calculations of Ref. 21.

It is interesting to note that for a fixed $\varepsilon_0/\varepsilon_f$, there is substantially less transverse flow developing in U + U collisions than in $p\bar{p}$ if hydrodynamic were to apply to $p\bar{p}$. However, an increasing initial energy density $\varepsilon_0(A) \propto R$ for large systems would still lead to a monotonic increase of $\langle p_\perp\rangle$ with A. In addition it is possible that the effective freezeout density ε_f also decreases with increasing R or A. Both effects tend to increase $\langle p_\perp\rangle$ with A slowly.

Applying the R = 1 fm curve in Fig. 3 to pp data[36], we still find a substantially larger increase of $\langle p_\perp\rangle$ with dN/dy than seen in the data. Furthermore pp and $\alpha\alpha$ data at ISR energies[37] do not show an increase of $\langle p_\perp\rangle$ with dN/dy. Most likely, finite mean free path effects ($R \sim \hbar/T$), non-scaling

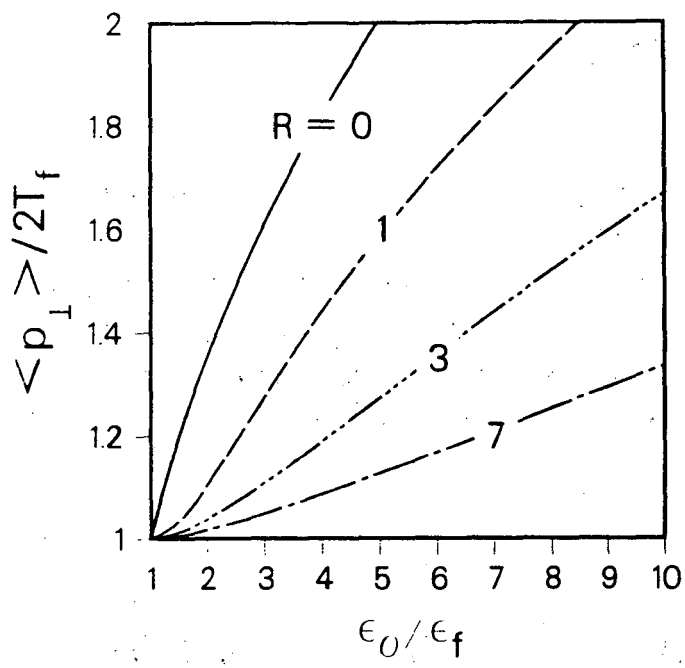

FIGURE 3
Enchancement of the average transverse momentum over that, $p_T = 2\ T_f$, due to the freezout temperature is shown as a function of the ratio of the initial to final energy densities. Reductions due to longitudinal expansion are shown for several size nuclei in fm.

of large fluctuation reactions, and large curvature and diffuseness affect $\langle p_\perp \rangle$ greatly in such small systems.

On the other hand, Van Hove has suggested[34] that the pp data may be reflecting the approximate first order character of the transition between the plasma and hadron phases. This can be seen by evaluating eq. (5.4) using a Bag model equation of state (solid curves in Fig. 1). We find[38] that

$$\langle p_\perp \rangle = p_T \begin{cases} (\epsilon_0/\epsilon_f)^{c_0/(1+c_0^2)} & : \ \epsilon_0 < \epsilon_H \\ (\epsilon_H/\epsilon_f)^{c_0/(1+c_0^2)} & : \ \epsilon_H < \epsilon_0 < \epsilon_Q \\ \left(\frac{\epsilon_H}{\epsilon_f}\right)^{\frac{c_0}{1+c_0^2}} \left(\frac{\epsilon_0 - B}{\epsilon_Q - B}\right)^{\frac{\sqrt{3}}{4}} & : \ \epsilon_Q < \epsilon_0 \ , \end{cases} \quad (5.15)$$

which shows the characteristic step structure anticipated in Ref. (34). As a first approximation to take longitudinal expansion into account the above result should be multiplied by $F(\epsilon_0, R)$.

Eq. (5.15) shows that the $\langle p_\perp \rangle$ barometer is most powerful when studied as a function of ϵ_0, i.e., dN/dy. In addition, because of the dependence of $\langle p_\perp \rangle$

on R due to $F(\epsilon_0,R)$ a systematic study of $\langle p_\perp \rangle$ vs. dN/dy and as a function of A is necessary to calibrate this barometer. Finally, we note that in the stopping regime, larger transverse flow can develop because high pressures are maintained for a relatively longer time by shock formation (see Fig. 2). In that case variation of $\langle p_\perp \rangle$ versus ϵ_0 is best achieved by varying the beam energy rather than studying fluctuations.

6. SEISMOMETERS

Up to now we have consider observables associated with average properties of the reaction. However, there may be large and interesting fluctuations around those averages. These fluctuations could arise as a result of fluctuations in the initial conditions[39] or as a result of violent processes[18,40], such as deflagrations or detonations, occurring in the plasma during the expansion phase. Of course, finite number effects always lead to fluctuations and only those fluctuations will be of interest that occur with a frequency greater than expected on trivial statistical grounds. In this section, I discuss a novel source of fluctuations that could arise if the transition between the plasma and hadronic worlds has a sharp step-like structure as illustrated in Fig. 1.

As shown by Baym et al.[21], for an equation of state as illustrated in Fig. 1, traverse hydrodynamic flow is unstable with respect to forming a shock wave. The precise conditions for shock formation were first discussed by Van Hove[40] and follow from combustion theory.[18] Consider a one dimensional quark-gluon plasma with energy density, pressure and flow velocity (ϵ_2,p_2,v_2) converting in a narrow region of width δ to hadronic matter with (ϵ_1,p_1,v_1). In the transition region, $\partial_\mu T^{\mu\nu} = 0$ holds. Integrating that equation across the transition region yields the unique values of flow velocities v_i as a function of ϵ_i,p_i.

$$v_1 v_2 = \frac{p_1 - p_2}{\epsilon_1 - \epsilon_2} , \qquad \frac{v_1}{v_2} = \frac{\epsilon_2 + p_1}{\epsilon_1 + p_2} . \tag{6.1}$$

By definition a deflagration (detonation) shock is one where $v_1 > v_2 (v_1 < v_2)$. In other words, the hadronic matter is accelerated (decelerated) with respect to the plasma in deflagration (detonation) processes.

While the hydrodynamic equations admit all solutions satisfying eq. (6.1), only those solutions are physical which lead to positive entropy production. In other words, the entropy current, su^μ, must satisfy[40] $s_1 u_1^z \geq$

$s_2 u_2^z$, where $s_i = \partial p_i / \partial T_i$ is the entropy density. For a bag model equation of state, the condition of positive entropy production gives[18]

$$\left(\frac{g_1}{g_2}\,\frac{\varepsilon_1}{\varepsilon_2 - 1}\right)^{1/2} \geq \frac{3\varepsilon_1 + \varepsilon_2 - 4}{3\varepsilon_2 + \varepsilon_1} \quad , \tag{6.2}$$

where g_i = number of boson + 7/8 x number of fermion helicity states in hadronic and quark-gluon matter respectively ($g_1/g_2 \sim 0.1$).

Numerical results[18,40] show that only deflagration shocks lead to normal hadronic matter densities $\varepsilon < \varepsilon_H$. Detonations require extreme supercooling $\varepsilon_2 \sim \varepsilon_Q/4$ and lead to extreme superheating $\varepsilon_1 \sim 20\ \varepsilon_H$ in the hadronic state. Therefore, detonations most likely do not occur. However, for $\varepsilon_2 \sim \varepsilon_Q$ deflagrations can lead to $\varepsilon_1 < \varepsilon_H$. The deflagration shock typically moves into the plasma at a very low velocity $v_{def} \sim 0.1$ while the ejected hadronic matter flows outward at rapidities on the order of the sound rapidity $y_s \cong 0.66$.

For a deflagration shock occurring on the (transverse) surface of the plasma this would yield an exhanced $p_\perp \cong p_T \exp(y_s) \sim 0.6$ GeV/c. Thus, rather high transverse momenta could arise from deflagration shocks if the plasma remained in a state with $\varepsilon_2 \sim \varepsilon_Q$. However, the ever present longitudinal expansion lowers ε_2 to $\varepsilon_Q/2$ on a time scale $\tau_o \sim 1$ fm/c. Simple deflagration shocks, however, cannot tolerate such supercooling.[18] Also given the small deflagration velocity, $v_{def} \sim 0.1$, their is a negligible change of the radius of the surface during the short time when a deflagration is allowed. Thus, the simple deflagration phenomena discussed in Ref. (40) is probably not applicable to the plasma produced in nuclear collisions. It remains an open question whether more complex deflagrations could arise on the surface. A hint of that comes from the analysis of deflagration bubbles.[18]

Figure 4 shows the development of a deflagration bubble in the plasma as a function of space and time. Detonation bubbles can also exist but as with

detonation shocks extreme supercooling of the plasma and extreme superheating of an expanding shell of hadronic matter is required. Growth of bubbles provides a mechanism to convert the cooling plasma into hadronic matter. The way in which a deflagration bubble comes about is through fluctuaions in the plasma in which a small domain of hadronic matter of radius r is formed. For r less than some critical radius, r_{crit}, the surface tension exceeds the gain in energy due to volume energy and the bubble collapses. For $r > r_{crit}$, the bubble is unstable against growth. Thus unlike surface deflagrations shocks, bubbles cannot be formed perturbatively. Indeed it was shown[21] that longitudinal hydrodynamic flow is stable with respect to small perturbations. Plasma at rest cannot turn into hadronic matter at rest by the

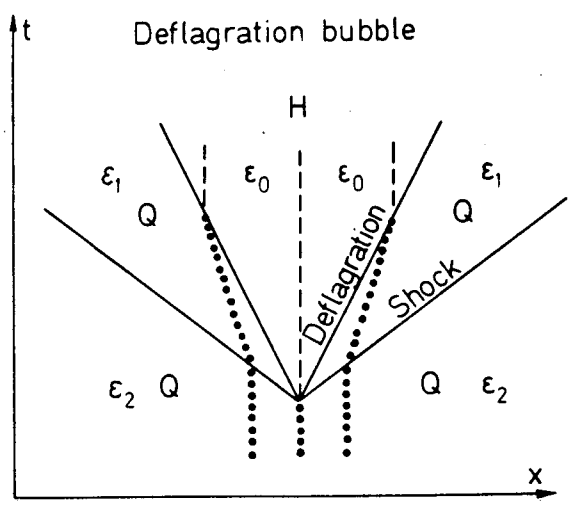

FIGURE 4
Space-time evolution of a deflagration bubble converting supercooled quark matter (Q) at rest into hadronic matter (H) at rest (dashed). Dotted line indicates quark paths.

requirement of entropy production. The trick which allows that transition is to send a preheating ordinary shock at high velocity outward from the bubble. That shock not only heats the plasma back up to $\varepsilon_1 \sim \varepsilon_Q$ but also accelerates the plasma outwards to a finite flow rapidity y_{f1}. This reheated moving plasma is then capable of deflagrating into hadronic matter at rest. Numerically, we found[18] that such bubbles can occur in supercooled plasma with $\varepsilon_Q > \varepsilon_2 > \varepsilon_Q/4$ and lead to $\varepsilon_1 \sim \varepsilon_Q$, $\varepsilon_o \sim \varepsilon_H$, and

deflagration and shock rapidities y_{def}, $y_{sh} \sim 1$. Therefore, they are ideal candidates for novel collective phenomena in the plasma relevant to nuclear collisions.

However, observable consequences of such explosive processes depend critically on the probability of forming an appropriate hadronic seeds with radius $r > r_{crit}$. That probability is exponentially sensitive to many uncertain parameters of the equation of state. Furthermore, as with familiar phase transition such as water boiling, impurities could be more important than tunneling processes in estimating that probability. In the absence of reliable estimates, we can only look for qualitative signatures of such phenomena.

We have considered three possible observables.[18] Consider a bubble seed with rapidity, y. Around that seed a deflagration bubble will grow in its rest frame as illustrated in Fig. 4. In order for the bubble to influence the plasma element with zero rapidity, the shock front must arrive at $z = 0$ before the plasma has broken up. That brake up time is roughly, $\tau_f \sim R$. Since the space-time position of a seed with rapidity y is ($t = \tau_0 \cosh y$, $z = \tau_0 \sinh y$) in the scaling regime and the front moves with near the velocity of light, only bubbles with rapidity $y < \ln(\tau_f/\tau_0)$ can influence the evolution of the plasma element with zero rapidity. Said another way, a bubble will be able to influence the evolution of the plasma in a total rapidity interval $\Delta y \sim 2 \ln R/\tau_0$ around its rest frame. Therefore, if a bubble is formed, then there will be medium range correlations with $\Delta y \sim 2$-4. In addition, since entropy is produced by the explosion, there will be an enhancement of dN/dy in a rapidity interval around the rapidity of the seed. A hint of unusual fluctuations of dN/dy has in fact been observed in a few cosmic ray events.[41] Thus, the study of fluctuations of dN/dy may provide a useful seismometer to search for explosive processes in nuclear collisions.

Another consequence of explosive bubbles is enhanced transverse momenta in an interval of Δy. In this case as the outward moving preheating and deflagration shocks encounter the transverse surface, the blast wave can give rise to a transverse flow of matter with rapidities, 1/2 - 1. This could show up as large $p_\perp > 1$ GeV/c of secondaries over a narrow interval of rapidity around the initial seed rapidity. Such anomalous fluctuations in $p_\perp$ versus y have been also seen in a few cosmic ray events.[41]

Finally, because of the azimuthal symmetry of the bubbles, the enhanced transverse momentum or, more precisely, high transverse energy density, $dE_\perp/dy$, would be associated with high circularity. Thus, explosive bubbles

would be easily distinguished from usual jet events where high $dE_\perp/dy$ is associated with small circularity.[42] Multiple medium jets would be the main source of background, which, however, could be estimated from pp data.

The observables discussed here require the study of multiparticle correlations. The fluctuations of dN/dy and $dE_\perp/dy$ and correlations between those fluctuations and global variables such as circularity, dC/dy, provide powerful handles in the search for such novel collective phenomena.

Finally, we note that the existence of deflagration bubble solutions discussed here[18] provides a clue to what may happen to surface deflagration.[40] When the interior plasma supercools, a single surface deflagration can no longer satisfy the hydrodynamic equations. However, a preheating imploding shock followed by an inward deflagration shock may exist under those conditions. Thus, the deflagration front probably breaks up into two fronts. This more complex surface deflagration can be thought of as turning the deflagration bubble inside-out. Details of such solutions are under investigation with P. Danielewicz.

7. SIGNATURES OR FORGERIES

In the quest for signatures of new phenomena there is always a danger of running into forgeries. We should demand in general positive identification via fingerprints. This metaphor, suggested by Cornelius Noack (private communications), rings especially true for studies involving nuclear collisions. We are after the simultaneous production and diagnostics of a speculative new form of matter--the quark-gluon plasma. However, the dynamical environment in which it is expected to be produced is very complex, as I have tried to emphasize. The pions, muons, and other products of the reaction would be observed whether or not this new state of matter is formed. It is their detailed spectra and correlations that could be influenced by the plasma phase. Unfortunately, no reliable background calculations exist yet with which those spectra and correlations can be calibrated. This is the most outstanding theoretical challenge in the near future. Nuclei are not macroscopic objects, though their length scales are several times larger than the scale, $\Lambda^{-1} \sim 1$ fm, relevant to hadronic processes. Finite number non-equilibrium effects will always contribute a non-negligable source of background.

One of the main lessons we have learned from the study of nuclear collisions at lower energies[43], $E_{lab} < 1$ GeV/A, is the necessity of treating those non-equilibrium processes explicitely. This required the

development of elaborate three dimensional Monte-Carlo cascade codes. Only now after a decade of analyzing data and refining the estimates of the non-equilibrium background are we beginning to have confidence that discrepancies between theory and experiment can be attributed to the interesting physics of dense, equilibrated nuclear matter. The debates are far from settled, but the equation of state of nuclear matter up to densities $\sim 4\rho_0$ seems now within reach.[44]

There is every reason to expect that the identification of the existence and properties of the plasma phase will be at least as difficult. Of course, we could get lucky, and long lived blobs of plasma or fractionally charged nuclei could provide the required fingerprints. However, in the absence of such stable or metastable exotic final fragments, only a long term dedicated effort involving the cross correlation of an arsenal of observables could lead to eventual success. No one experiment is likely to be decisive. Elaborate devices are called for that can measure the exclusive characteristics of such reactions on an event by event basis. The most useless of all measurements are untriggered inclusive yields. Those inclusive yields are easiest to forge theoretically. Indeed, as with the 1 GeV/A domain there already exist a zoo of models capable of fitting every inclusive data point. The only hope of weeding out models is through more exclusive data or inclusive data triggered on exclusive characteristics (multiplicity, circularity, transverse energy, etc.).

In this talk, I have considered only three generic observables. This was to illustrate a few of the current ideas and problems that make this new field so exciting. For a discussion of other observables associated with strangeness see the talk of J. Rafelski elsewhere in these proceedings. Some of the topics not discussed but which also warrant serious consideration are (1) the formation of exotic objects associated with color and quark degrees of freedom,[45] (2) $\pi\pi$, KK, multi π, and speckle interferometry[1] to uncover space-time characteristics of the reaction, (3) observables, such as resonance widths, that could probe the chiral symmetry restoration transition,[46] (4) correlations between spect ator fragments and participant fragmentation yields as a measure of transport properties,[29] (5) light nuclei (d, t, α) in the fragmentation regions as a measure of entropy production,[43] (6) jet production and quenching as a measure of stopping power in the plasma,[47] and (7) the stopping power of nuclei in the 10 GeV/A region[48-49] which is crucial for the producton of high baryon density plasmas.[25]

Sorting the signatures from the forgeries and hunting for fingerprints of

the quark-gluon plasma promise to be one of the most fun detective stories in the coming decades.

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