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PROBES OF THE QUARK-GLUON PLASMA AS IT MIGHT BE PRODUCED IN ULTRA-RELATIVISTIC NUCLEAR COLLISIONS*

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The energy densities which might be achieved in ultrarelativistic nuclear collisions are discussed. Using these estimates, promising probes of a quark-gluon plasma as it might be produced in such collisions are reviewed. I discuss in detail the emission of photons and di-leptons. The consequences of hydrodynamic expansion and a first order phase transition are explored for the transverse momentum spectrum of hadrons. Fluctuations in the rapidity distribution of hadrons are also discussed as a possible signal for a first order phase transition. The possibility that copious production of strange particles may signal the production of a quark-gluon plasma is critically assessed.

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I shall discuss the possible experimental probes of the quark-gluon plasma as it might be produced in ultra-relativistic nuclear collisions. I shall concentrate on the central region of collisions of large nuclei, $A \gtrsim 200$, for head-on collisions at extremely high energies, $E_{CM}/A > 50$ GeV/Nucleon. A picture of such a collision is shown in Fig. 1 [1–2]. At some time τ_0 after the two nuclei pass through one another, matter begins to form between them. In the inside-outside cascade picture of the collision, this forming matter is assumed to be non-interacting until after the time τ_0. The rapidity of the particles which constitute newly forming matter is therefore

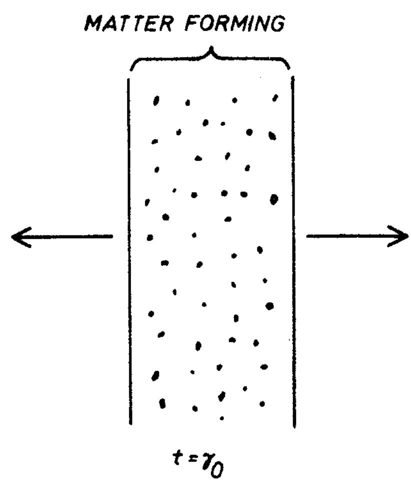

Fig. 1. A nucleus-nucleus collision

$$y = \tfrac{1}{2} \ln \left\{\frac{1+v}{1-v}\right\} = \tfrac{1}{2} \ln \left\{\frac{t+x}{t-x}\right\}, \tag{1}$$

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where v is a particle velocity, t is the time measured from the initial time of the collision, and x is the longitudinal coordinate measured from the position of the collision. This correlation between momentum and space-time persists after the time τ_0 as a consequence of the hydrodynamic equations, and may be taken to be valid for all times.

The energy density of matter at the formation time τ_0 is [1]

$$\varrho = \frac{1}{\pi A^{2/3}} \frac{dN}{dx} \langle m_t \rangle = \frac{1}{\pi A^{2/3}} \frac{dN}{dy} \frac{\langle m_t \rangle}{\tau_0}. \tag{2}$$

This result has been used to estimate the energy densities achieved in the ultra-relativistic collisions observed by the JACEE cosmic ray experiment [3]. If $\langle m_t \rangle \sim 0.4$ GeV and $\tau_0 \sim 1$ fm/c, the dN/dy distributions observed for intermediate A nuclei extrapolated to heavy nuclei such as uranium predict energy densities of 5–10 GeV/fm^3. Such energy densities may be sufficient to produce a quark-gluon plasma [4].

Recent results for hadron-nucleus collisions indicate that the formation time $\tau_0 \sim 1$ fm/c may be a little large [5]. The dependence of the energy density of Eq. (2) upon τ_0 is not trivial. By the uncertainty principle,

$$\langle m_t \rangle \gtrsim 1/\tau_0 \tag{3}$$

so that

$$\varrho \gtrsim \frac{1}{\pi A^{2/3}} \frac{dN}{dy} \frac{1}{\tau_0^2}. \tag{4}$$

Since the energy density of a quark-gluon plasma scales as T^4 where T is the temperature, and the maximum achieved temperature is $T \sim \tau_0^{-1/2}$.

In a nice analysis presented by D. Kisielewska, the possible values of τ_0 are extracted from hadron-nucleus and lepton-nucleus experimental data [6]. The range of values consistent with these data are determined to be $1/5 < \tau_0 < 1$ fm/c. The preferred values are 1/2–1/3 fm/c. (It should be noted that in string models of nucleus-nucleus collisions, the formation time depends upon A, and may be considerably smaller for large A nuclei than is the case for hadron-nucleus collisions) [7]. If we consider a range $1/20 < \tau_0 < 1$ fm/c, the corresponding energy densities and temperatures are $\varrho \sim$ 5–5000 GeV/fm^3 and $T \sim$ 0.2–1 GeV.

Another method of estimating the energy densities achieved in ultra-relativistic nuclear collisions has been advocated by Matsui and Gyulassy [8]. They make use of the observation that the rapidity density $\frac{dN}{dy}$ is conserved in isentropic expansion. The hydrodynamic equations may be integrated backwards from the final time at which the system breaks up to the initial time τ. They derive [8]

$$\varrho \sim 1.6 \left\{ \frac{1}{\tau_0 \pi R^2} \frac{dN}{dy} \right\}^{4/3}. \tag{5}$$

The difference between this estimate and that of Eq. (4) is that the *dN/dy* distribution quoted here is that observed in the final state, whereas in Eq. (4) it is only the initial rapidity density which appears, and this might be changed if the system produces entropy as it expands. Usually it is assumed that the system thermalizes at formation, that is, the formation time is of the order of the collision time and that the initial distributions are close to thermal distributions. In this circumstance, it should be a fair approximation to treat the expansion as isentropic, and the initial rapidity density may be identified with the final one. In general, this cannot be true since Eqs. (4)–(5) do not agree. For practical purposes, the agreement is quite good however. If the formation time varies between 0.1 and 1 fm, the estimated energy density varies by two orders of magnitude, where the difference between the two different estimates varies only by a factor 3–5. As order of magnitude estimates, either relation is acceptable.

If one requires that the system instantly thermalize, then both Eqs. (4) and (5) must be valid. This can only be true if the formation time is A dependent. An estimate of this dependence is

$$\tau \sim 0.4\, A^{-1/6} - 0.6\, A^{-1/3}. \tag{6}$$

For heavy nuclei, formation times of 0.1–0.2 fm would be consistent with either relationship.

If the formation time is small and if the formation time is of the order of the collision time, then the condition that the system be thermalized and expand according to the equations of perfect fluid hydrodynamics, that is, non-viscous hydrodynamics, seems on much firmer ground than is the case for larger formation times. Notice that according to the hydrodynamic equations, the initial time and temperature and the final time and temperature are related as

$$\tau_f = \tau_i \left\{ \frac{T_i}{T_f} \right\}^3 . \tag{7}$$

If the initial temperature is 250 MeV, the time the system takes to cool to a temperature of 150 MeV is only a factor of five times the initial time. If the initial temperature is a factor of two larger, this ratio increases by nearly an order of magnitude, and hydrodynamic methods are probably on a somewhat better foundation.

Since the width of the fragmentation region is given by

$$y_{\text{frag}} \sim \ln R_{\text{nuc}}/\tau_0 \tag{8}$$

experimental measurements of the width of this region may aid in a resolution of τ_0. Such a measurement might be to determine the rapidity distribution of baryons minus anti-baryons, Fig. 2, or $\pi^+ - \pi^-$ mesons, Fig. 3. The values of E_{CM}/A required to produce a baryon free central region depend upon τ_0, and for the values of τ_0 above are 15–300 GeV/Nucleon.

The value of τ_0 needed in the above analysis might be determined by using the Hanburry-Brown-Twiss effect [9–10]. The case that these measurements are useful for this purpose has not been made since hadronic final state interactions probably obscure a study of the formation process. Pion interferometry is probably most useful for studying the space-time dynamics of the hadronization process, or the break-up of the system as it freezes out of thermal equilibrium [9–10]. Again, there has been little theoretical analysis of this problem.

The formation time τ_0 might be measured and the validity of the formation time idea might be tested in hadron-nucleus collisions. As mentioned above, data analyzed by Busza and Goldhaber suggest that at available fixed target energies, projectile hadrons are stopped much more efficiently than might be expected from an inside-outside cascade picture with a formation time $\tau_0 \sim 1$ fm [5]. The Kisielewska analysis is consistent with an inside-outside cascade picture if formation times somewhat smaller than 1 fm are assumed [6]. Also, a new method of analysis has been developed by Białas to measure the formation time for hadrons, a time which should be longer than the matter formation

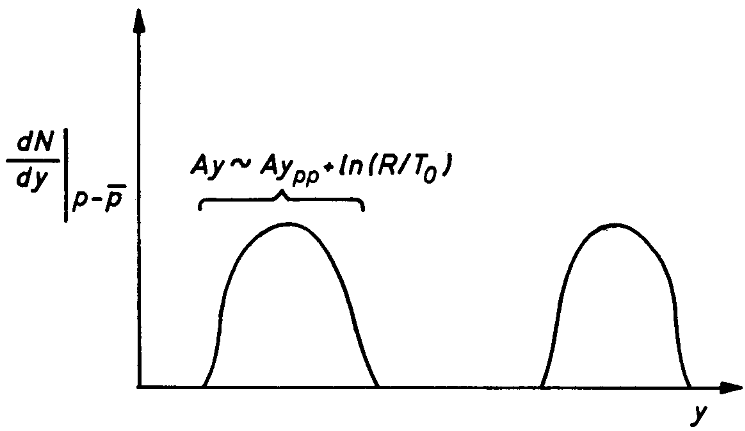

Fig. 2. dN/dy for proton minus anti-proton

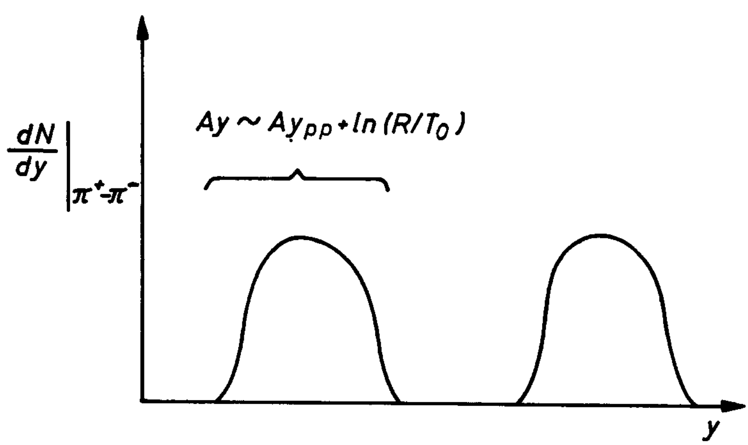

Fig. 3. dN/dy for π^+ minus π^-

time for quarks and gluons [11]. When baryon production is analyzed, a baryon formation time of 1.5 ± 0.5 fm results.

There is at present no consensus on methods for analyzing data from hadron-nucleus collisons and extracting a matter formation time. A very promising idea has been put forward by Hwa who suggests that the deviations of hadron-nucleus scattering from hadron-hadron in the hadron fragmentation region may cleanly isolate the effects of matter formation [12]. The argument is that within the context of an inside-outside cascade model, the probability that an inelastically produced particle forms inside a target nucleus is small if the particle is very energetic. The probability of rescattering is $P \sim e^{-x/\lambda}$ where λ is the mean free path for rescattering. This mean free path is roughly the distance it takes a particle to form, since upon formation in the target, it has a large probability to rescatter, so that

$$\lambda \sim \frac{E\tau_0}{\langle m_t \rangle}. \tag{9}$$

Since the particle has only a small chance to rescatter, its modification to hadron-hadron distributions, δF, should be small and computable in a single rescattering approximation,

$$\delta F \sim \frac{A^{1/3}}{\lambda}.$$

A detailed, but still controversial analysis by Zahir and Hwa gives a λ which is consistent with the time dilation expected in the inside-outside cascade models [12].

On the other hand, Glauber theory type models such as those employed by Kapusta and Csernai, and by Wong also seem to fit some of the hadron-nucleus data, and such models would be inconsistent with inside-outside cascade models unless the formation time was very small [13–14].

A proper resolution of these theoretical models may require more experimental data. Experiments with $E_{\rm lab} \sim 10$–100 GeV, for a wide range of A, including hadron--proton and hadron-deuterium, would be useful. A wide range of x_F is probably best for the type of analysis suggested by Hwa [13], but if good cascade models of hadronic interactions are developed along the lines of Kapusta and Csernai and of Wong, then coverage of the central and target fragmentation regions are also necessary [13–14]. The data on targets of various A including protons and deuterium should come from the same experiment to sort out systematic experimental biases.

The study of the time development of the hadronic matter distribution of matter as it is produced in ultra-relativistic heavy ion collisions has been initiated by the hydrodynamic computations of Bjorken [1], by Kajantie, Raitio, and Ruuskanen [15]. In these computations, the longitudinal expansion of the matter is studied. In the computations of Kajantie et al., the fragmentation region as well as the central region is simulated. In later computations by Baym et al. and Białas et al., the transverse expansion is also computed [16–17]. The hydrodynamic treatment concludes that for formation time $\tau_0 \sim 1$ fm, energy densities of 2–10 GeV/fm^3 may be obtained in the central region, and 0–2 GeV/fm^3 at various rapidities of the fragmentation region, with smallest values at the largest values of Feynman x. The compression of the baryon number density in the fragmentation region is 0–2 times that of ordinary nuclear matter. The transverse expansion calculations have been done only for the central region. Such expansion consists of transverse rarefraction and takes place over time scales large compared to longitudinal expansion for large $A \gtrsim 200$ nuclei.

In order for the hydrodynamic treatment to be valid, the mean free paths for quarks and gluons must be small compared to the spatial dimensions of the matter produced in the collision. Detailed computations seem to verify such a treatment for large A nuclei [18–19]. A much more stringent test of the validity of perfect fluid hydrodynamics is that the collision time be small compared to the longitudinal expansion time. Estimates of the collision time for appropriate energy densities are $\tau_c \sim 0.1$–1 fm. Since a hydrodynamic treatment is valid if $\tau \gtrsim \tau_c$, the intrinsic error in these computations does not allow a good resolution of the time after which a hydrodynamic computation is reasonable. At the very least, viscous corrections to the perfect fluid hydrodynamic equations are probably important at early times in the hydrodynamic expansion, and there is certainly substantial entropy production for $\tau \sim \tau_c$.

Some measure of the degree of thermalization and the validity of a hydrodynamic treatment have been suggested by Shuryak [20]. In the transverse expansion of the matter produced in nuclear collisions, a transverse flow velocity develops. Particles of all masses flow with the velocity of the fluid. The more massive particles therefore have their transverse momentum enhanced more relative to lighter particles. A detailed treatment of this problem using the methods developed in Refs. [16–17] would be useful.

If there is a first order phase transition in hadronic matter, the transverse momentum distribution of hadrons may be drastically altered. Following Shuryak, the transverse momentum distributions receive a contribution due to transverse hydrodynamic expansion, and a thermal contribution due to the breakup of the system at some temperature T [20–21]

$$\langle p_t \rangle = \langle p_t \rangle_{\mathrm{hydro}} + \langle p_t \rangle_{\mathrm{thermal}}. \tag{10}$$

The hydrodynamic contribution arises from work which is done on constituents of the matter as the matter is driven into the vacuum. This work is produced from the pressure difference of the matter and the vacuum. Consider the work done as a function of the energy density achieved by the matter in a nuclear collision. As the energy density increases, the pressure will increase except when the energy density is in the region of a mixed phase of hadronic matter and quark-gluon plasma. For such energy densities, the pressure remains constant, as is shown in Fig. 4a. As the pressure increases, the transverse momentum due to hydrodynamic expansion increases. In the region of the phase transition, the pressure remains constant, as does the transverse momentum of hadrons.

This general feature of the hydrodynamic expansion coupled with a phase transition may be explored by plotting measured transverse achieved energy densities, as is shown in Fig. 4b [20–21].

Such a plot has been made by the JACEE cosmic ray collaboration, and is shown in Fig. 5 [22]. Although the inferred energy density is somewhat model dependent, the sharp break in the transverse momentum distribution is quite suggestive.

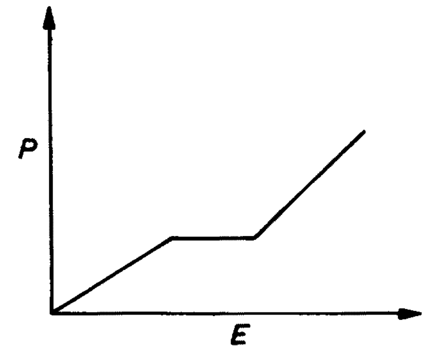

Fig. 4a. P vs E

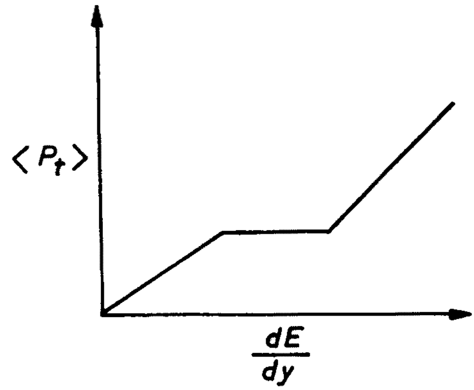

Fig. 4b. $\langle p_t \rangle$ vs dE/dy

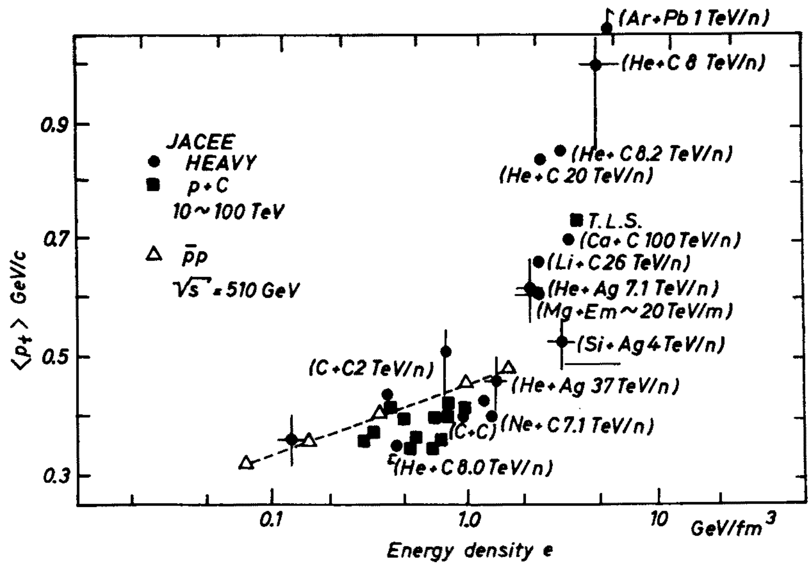

Fig. 5. $\langle p_t \rangle$ vs dE/dy as measured by JACEE experiment

Studies by Heinz and by Białas and Czyż suggest that color plasma oscillations may play an important role in non-equilibrium processes early in the expansion of matter produced in ultra-relativistic nuclear collisions [23–24]. A typical color oscillation is shown in Fig. 6. In such an oscillation, the local color charge density is analogous to the electromagnetic charge density in the electromagnetic plasma oscillation. For the oscillation shown in Fig. 6, both a color and electromagnetic oscillation is set up, and soft electromagnetic radiation may be emitted from the oscillating charge density. Białas and Czyż estimate that for a not unreasonable spectrum of color plasma oscillations which may be characteristic of nuclear collisions,

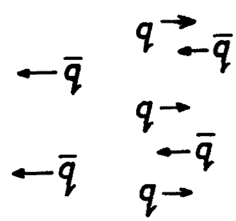

Fig. 6. A color plasma oscillation

a substantial fraction of the hadronic energy may be emitted in low p_t photon radiation [24].

Electromagnetic probes of the nuclear collisions may be characterized by the value of the transverse mass, $M_t = \{p_t^2 + M^2\}^{1/2}$. For di-lepton and photon transverse masses $M_t \lesssim R_{\rm nuc}^{-1} \sim 30$ MeV for Uranium, the photons and di-leptons are coherently produced [25–26]. These particles may be copiously produced in the nuclear fragmentation regions, or in the central region by fluctuations in the charge distributions of mesons. Detailed measurements of these distributions in correlation with measurements of the charge distributions, dN/dy, may probe the electromagnetic plasma oscillation, since as a consequence of this oscillation, radiation with frequency less than the electromagnetic plasma frequency is strongly absorbed. The detailed computation of this absorption is complicated by the finite size of the matter produced in nuclear collisions, and since the matter density is a decreasing function of time, the plasma frequency is time dependent.

In the central region, these coherently produced low transverse mass photons are generated by charge fluctuations. In order to have a large number of photons, a large charge fluctuation must be generated. At sufficiently small p_t, however, the coherently produced photons will dominate over those arising from hadronic decays. In the fragmentation region, the net charge carried by the projectile nucleus generates a large number of low p_t photons [25–26]. The total number is $N_{\rm tot} \sim Z^2\alpha \ln \dfrac{\omega_{\max}}{\omega_{\min}}$, where Z is the total number of struck nucleons, and $\omega_{\max}$ and $\omega_{\min}$ are maximum and minimum observed frequencies. These limiting frequencies are determined by detectors and backgrounds. Typically the maximum frequency is limited by π^0 decay photons and is $\omega \sim 15$ MeV in the rest frame of the struck nucleus. The low frequency cutoff is more difficult to estimate. The sensitivity of this result to Z suggests that measurements of these low p_t photons may provide a good impact parameter meter. The number and distribution of these emitted photons may also be used to infer the rapidity distributions of charged particles. The total number of photons emitted in a reasonable frequency range for reasonable estimates of the rapidity distributions of charged particles in head-on collisions of large A nuclei is $N_{\rm tot} \sim 50-500$. Photons emitted from a beam projectile nucleus, or from colliding nuclei are in a small angular region $\Delta\theta \sim 1/\gamma$ where γ is the Lorentz γ factor of the nucleus. The energy of these photons is $E \sim \gamma\omega$. For a 100 GeV beam, a reasonable range of these parameters is $\Delta\theta < 0.6°$ and $E < 1.5$ GeV.

For larger values of the transverse mass, photons and di-leptons may be approximated as elementary probes with mean free paths large compared to the size of the matter produced in a nuclear collision. Such probes have advantages over hadrons, since hadrons strongly interact and their distribution is characteristic of matter either at the surface or at late times when the matter is at such a low density that hadrons cease interacting. Photons and di-leptons probe the matter at early times when it is hot and dense [27].

A systematic study of photon and di-lepton production may start with a study of the thermal expectation value of the electromagnetic current-current correlation function [27–29]

$$W^{\mu\nu}(q) = \int d^4x e^{iqx} \langle J^\mu(x) J^\nu(0) \rangle. \tag{11}$$

The rate for thermal emission is related to this structure function as

$$Rate/\text{Volume} \sim e^2 L^{\mu\nu}(q) W_{\mu\nu}(q), \tag{12}$$

where $L^{\mu\nu}$ is a computable lepton polarization tensor.

The properties of $W^{\mu\nu}$ are formally very similar to that of $W^{\mu\nu}$ for deep inelastic scattering of leptons from hadrons. The only difference is that there is a thermal expectation value here and not the matrix element between proton states, and that the photon momentum is either timelike or lightlike, not spacelike. The fluid four velocity u^μ is a timelike vector analogous to the proton four momentum of deep inelastic scattering. Because of these formal similarities, $W^{\mu\nu}$ may be written in terms of invariant structure functions A and B as

$$W^{\mu\nu} = \{q^2 g^{\mu\nu} - q^\mu q^\nu\} A(q^2, u \cdot q, T, \Lambda) + \{g^{\mu\nu}(u \cdot q)^2 - (u^\mu q^\nu + u^\nu q^\mu) u \cdot q + u^\mu u^\nu q^2\} B(q^2, u \cdot q, T, \Lambda). \tag{13}$$

The QCD scale parameter Λ is written explicitly in this equation.

The properties of this structure for large q may be studied and scaling behaviour analogous to that for deep inelastic scattering functions may be extracted. At very high temperatures, $T/\Lambda \gg 1$, these structure functions may be computed in perturbation theory. In the large q limit, a detailed analysis of the asymptotic behaviour of A and B gives

$$A \to e^{\beta u \cdot q} \bar{A}((u \cdot q)^2/q^2, \beta\Lambda, q^2/\Lambda^2), \tag{14}$$

$$B \to e^{\beta u \cdot q} \bar{B}((u \cdot q)^2/q^2, \beta\Lambda, q^2/\Lambda^2). \tag{15}$$

Here $\beta = 1/T$. In the limit of large temperatures, the structure function B vanishes. The asymptotic scaling property of A and B shown in Eqs (14)–(15) is called thermal scaling. The Boltzmann weight factor $e^{\beta u \cdot q}$ and the dependence upon $(u \cdot q)^2/q^2$ are different than that of $W_{1,2}$, for deep inelastic scattering, which has no weight factor, and is a function of the Bjorken x variable, $x = q^2/(p \cdot q)$.

The structure function $W^{\mu\nu}$ must be folded into the hydrodynamic equations before experimental distributions of photons and di-leptons may be computed. I shall only state the results of such an analysis here. For photon and di-lepton transverse masses large compared to the temperature, the photon and di-lepton rapidities and transverse masses are closely correlated to the rapidity of the plasma from which they were emitted and the temperature of the plasma at the emission time. Direct computation gives

$$y_{\text{photon}} \sim y_{\text{plasma}}, \tag{16}$$

$$M_t \sim \left(\frac{2}{v_s^2} + \tfrac{1}{2}\right) T, \tag{17}$$

where v_s is the sound velocity of the matter. For sound velocities characteristic of an ideal quark-gluon plasma,

$$M_t \sim 6.5\ T \tag{18}$$

and the assumption that the transverse masses are large compared to the temperature seems a posteriori justified. For temperatures of $T \sim 100-500$ MeV, transverse masses of $M_t \sim 0.6-3$ GeV are emitted with greatest strength.

The absolute rate for photon and di-lepton emission is extremely sensitive to the maximum temperature achieved in the collision and to the sound velocity of the hadronic matter. Roughly three orders of magnitude result from each of these uncertainties. Probably, the rate for di-lepton production is 1–10^4 orders of magnitude times the Drell-Yan rate extrapolated to this mass region. The uncertainty in the total rate is reflected in the A dependences of the total rates. A measurement of the A dependence of the total emission rate provides a check on dynamical assumptions used in computing the thermal emission spectrum.

The shape of the emission spectrum of di-leptons is quite different from that of Drell-Yan. The distribution is a pure power of M_t if the plasma acquires a large enough temperature,

$$\frac{dN}{dM^2 dy d^2 p_t} \sim \left\{\frac{1}{M_t}\right\}^{2/v^2_s}. \tag{19}$$

As a consequence, the transverse momentum and the mass of the di-lepton pairs are strongly correlated. The structure function for di-lepton emission from a high temperature plasma, corrected for expansion involves only one structure function,

$$\Omega^{\mu\nu} \sim \{q^2 g^{\mu\nu} - q^\mu q^\nu\}\Omega_1. \tag{20}$$

As a consequence of the correlation between temperature and transverse mass, any discontinuity as a function of temperature also appears as a function of transverse mass. If there is a first order phase transition, the electromagnetic current-current correlation function should be discontinuous across the phase transition, and this may appear as a discontinuity in the transverse mass spectrum as shown in Fig. 7. If the system exists in a mixed phase for a long time, the thermal emission spectrum will attain a contribution of $e^{-M_t/T_{\mathrm{p.t.}}}$, since the mixture of plasma and hadronic matter will emit at the phase transition temperature $T_{\mathrm{p.t.}}$.

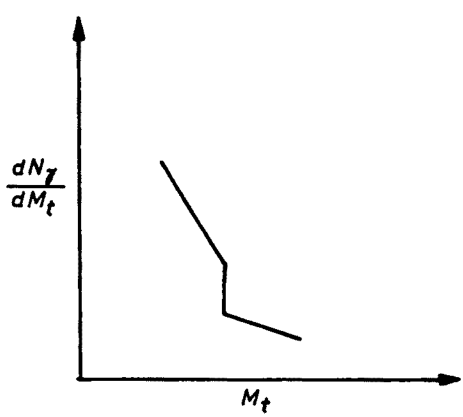

Fig. 7. A possible transverse photon mass distribution

The physics which may be studied by photons and di-leptons is characterized by transverse mass values. The range of $M_t \sim 0.6-3$ GeV may be dominated by thermal emission. For very low M_t values, coherent emission processes dominate. For masses 30 MeV $< M <$ 200 MeV, the effects of coherence begin to subside, and incoherent production processes begin to dominate. In the region $200 \lesssim M_t \lesssim 600$ MeV, incoherent processes should dominate. In these low mass regions it might be possible to probe the effects of chiral symmetry restoration [30]. The rate of production of such low transverse mass particles might be very sensitive to the constituent quark masses, since the di-lepton production process shown in Fig. 8 vanishes below the threshold $q^2 < 4m_q^2$. The outstanding problem in this low mass region is resolving background processes arising from hadronic decays. A thorough theoretical analysis of backgrounds and a comparison with emissions from a quark-gluon plasma has not yet been carried out.

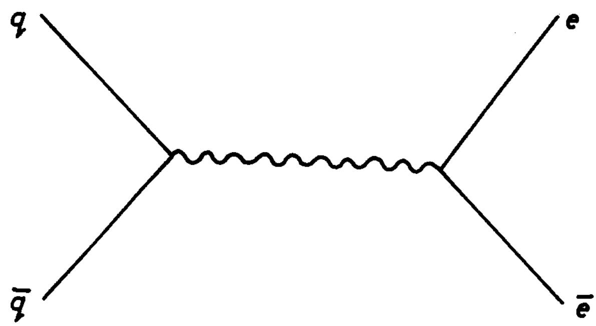

Fig. 8. Quark annihilation into lepton pairs

At large transverse masses $1\text{–}5 \lesssim M_t \lesssim 10\text{–}20$ GeV, there should be corrections to the Drell-Yan emission rates arising from the pre-equilibrium distributions of quarks and gluons. At present a theory of these distributions is lacking, but the development of such a theory is necessary to put the production of a quark-gluon plasma in ultra-relativistic nuclear collisions on a stable foundation.

At transverse masses $M_t \gtrsim 10\text{–}20$ GeV, the Drell-Yan process should dominate.

The production of strange particles has long been suggested as a signal for the production of a plasma [31]. The ratio of strange to non-strange anti-baryons might retain some trace of an abundance of strange quarks and anti-quarks produced in a plasma. This conclusion is on somewhat shaky ground since in the hydrodynamic expansion of the plasma, the strange quarks and anti-quarks may become diluted. Also, a recent computation of Redlich suggests that the abundance of strangeness in a hadronic gas may not be so far different from that of a quark-gluon plasma [32]. A proper theoretical assesment of strangeness production probably needs non-perturbative input from lattice Monte-Carlo computations, and a thorough analysis of the effects of hydrodynamic expansion.

Charm particle production may also be important if sufficiently high plasma temperatures are achieved, $T \gtrsim 500$ MeV [33]. Corrections due to hydrodynamic expansion are probably less important for charmed particles than for strange particles since the charmed quark hadronic cross section is small $\sigma < 1$ mb.

An extremely speculative experimental probe of quark-gluon plasma production may be in multi-particle correlations, and in large scale rapidity fluctuations. Such correlations and fluctuations may arise as the matter participating in a nuclear collision tries to negotiate a first order transition [34–36]. A variety of scenarios are possible all of which involve the production of large scale density fluctuations over rapidity intervals $\Delta y \gtrsim 1$. In the collisions of heavy nuclei, such a rapidity interval may include several hundred to several thousand particles, and large scale fluctuations should be separable from statistical fluctuations. These density fluctuations may be generated by superheating, supercooling or the spinodal decomposition of the plasma. They might occur in baryon, anti-baryon or meson distributions [37]. There might also be p_t enhancements if the density fluctuations are accompanied by burning or explosive phenomenon. Backgrounds such as jet production may be ruled out by the azimuthal angle distributions.

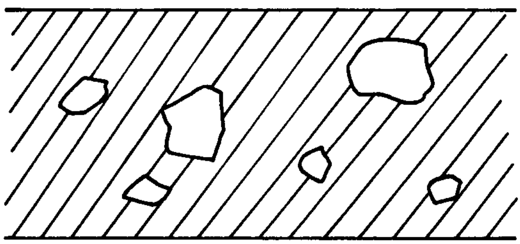

Fig. 9. Large scale density fluctuations in the plasma

The restoration of chiral symmetry may have striking consequences for the widths and masses of resonances produced in a quark-gluon plasma [38]. Since the chiral symmetry transition is expected to be abrupt, the masses and resonances may abruptly change their character as the energy density achieved in a collision increases. Correlations of the type described for $\langle p_t \rangle$ vs dN/dy may be useful here. Also, anomalously large baryon or anti-baryon production might accompany chiral symmetry restoration. As a chiral symmetric world cools through a first order phase transition, light mass baryons in the chirally symmetric phase might become clustered into plasma droplets, as shown in Fig. 9. The large scale density fluctuations which characterize first order phase transitions might appear as large scale fluctuations in the rapidity distribution of baryons and anti-baryons.

The distribution of jets produced in ultra-relativistic nuclear collisions may provide probes of the space-time evolution of the plasma, and the matter distribution produced in the collision. For example, the occurrence of single jets, where one jet has been absorbed as it passes through the plasma, shown in Fig. 10, provides a measure of the mean free path of quarks and gluons in hadronic matter [39].

Various speculations concerning the existence of exotic stable or metastable forms of matter have been suggested. Examples are Lee-Wick matter [40], stable or metastable droplets of chirally symmetric strange matter [41–43] or charmed matter [44] and even metastable Higgs meson matter [45]. Also, arguments have been proposed that free quarks might be easier to produce in nuclear collisions than in e^+e^- or pp collisions [46]. The precise nature and the probability that such matter exists are difficult to determine, but the revolutionary character of its discovery justifies a generic search. Hints that such new forms of matter exist in cosmic rays have long been suggested by mountain-top emulsion chamber experiments [47]. Such new forms of matter fortunately have generic characteristics which distin-

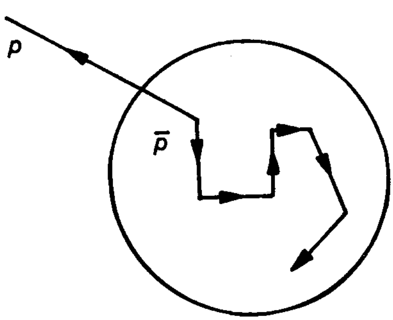

Fig. 10. Single jet production

guish it from ordinary nuclear matter. For free quarks, the charge is a signature. For particles whose production is associated with a conserved quantum number, the exotic particles should be produced in pairs. In general, the charge to mass ratios of exotic particles should bear no simple relation to that of ordinary nuclei. The penetrating power and cross sections are in general different for ordinary particles with the same Q or Q/A. If an exotic particle decays, the multiplicity might be anomalously large, and the p_t distribution might not be typical of either a nuclear break-up or a hadronic interaction. Since the exotic particle may carry charge, strangeness, or baryon number, the flavor composition and charge to neutral composition of the final state may be anomalous. Secondaries of the decay may also themselves have anomalous interactions or decays.

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REFERENCES

[1] J. D. Bjorken, *Phys. Rev.* **D27**, 140 (1983).
[2] K. Kajantie, L. McLerran, *Nucl. Phys.* **B214**, 261 (1983); *Phys. Lett.* **119B**, 203 (1982).
[3] T. Burnett et al., *Phys. Rev. Lett.* **50**, 2062 (1983).
[4] See Proceedings of Quark Matter 84, Helsinki, Finland 1984.
[5] W. Busza, A. Goldhaber, *Phys. Lett.* **139B**, 235 (1984).
[6] D. Kisielewska, *Acta Phys. Pol.* **B15**, 1111 (1984).
[7] H. Ehtamo, J. Lindfors, L. McLerran, *Z. Phys.* **C18**, 341 (1983).
[8] M. Gyulassy, T. Matsui, *Phys. Rev.* **D29**, 419 (1984).
[9] A. Goldhaber, *Nature* **275**, 114 (1978).
[10] M. Gyulassy, S. K. Kauffmann, L. W. Wilson, *Phys. Rev.* **C20**, 2267 (1979).
[11] A. Białas, Max Planck Institute Preprint, MPI-PEA-PTh-33/84 (1984).
[12] R. Hwa, M. S. Zahir, Univ. of Oregon Preprint, OITS-256, April 1984.
[13] R. Hwa, *Phys. Rev. Lett.* **52**, 492 (1984).
[14] J. Kapusta, *Phys. Rev.* **C27**, 2037 (1983); L. Csernai, J. Kapusta, Univ. of Minnesota Preprint, UMNTNP-70/1984 (1984); C. Y. Wong, *Phys. Rev. Lett.* **52**, 1393 (1984).
[15] K. Kajantie, R. Raitio, P. V. Ruuskanen, *Nucl. Phys.* **B222**, 152 (1983).
[16] A. Białas, W. Czyż, *Acta Phys. Pol.* **B15**, 229 (1984).
[17] G. Baym, B. L. Friman, J. P. Blaizot, M. Soyeur, W. Czyż, *Nucl. Phys.* **A407**, 541 (1983).
[18] A. Hosoya, K. Kajantie, Helsinki University Preprint, HU-TFT-83-62 (1983).
[19] P. Danielowicz, M. Gyulassy, LBL Preprint, LBL-17278 (1984).
[20] E. V. Shuryak, *Phys. Rep.* **61**, 71 (1980); E. V. Shuryak, O. V. Zhirov, *Phys. Lett.* **89B**, (1979).
[21] L. van Hove, *Phys. Lett.* **118B**, 138 (1982).
[22] O. Miyamura, Proceedings of Quark Matter 84, Helsinki, Finland 1984.
[23] U. Heinz, *Phys. Rev. Lett.* **51**, 351 (1983).
[24] A. Białas, W. Czyż, *Phys. Rev.* **D30**, 2371 (1984).
[25] J. D. Bjorken, L. McLerran, Fermilab-Pub-84/35-T.
[26] J. Kapusta, *Phys. Rev.* **C15**, 1580 (1977).
[27] E. L. Feinberg, *Nuovo Cimento* **34A**, 391 (1976).
[28] E. V. Shuryak, *Phys. Lett.* **78B**, 150 (1978); *Sov. J. Phys.* **28**, 408 (1978).
[29] L. McLerran, T. Toimela, Fermilab-Pub-84/84-T.
[30] J. Kapusta, *Phys. Lett.* **136B**, 201 (1984).
[31] J. Rafelski, B. Muller, *Phys. Rev. Lett.* **48**, 1066 (1982); T. S. Biro, B. Lukacs, J. Zimanyi, H. Barz, *Nucl. Phys.* **A386**, 617 (1982).
[32] K. Redlich, Technische Hochschule Darmstadt Preprint, IKDA 84/12.

[33] J. Cleymans, *Phys. Lett.* **127B**, 375 (1978); J. Cleymans, R. Philippe, *Z. Phys.* **C22**, 271 (1984); J. Cleymans, G. Vanderzande, Bielefeld Preprint, BI-TP-84/04 (1984).
[34] L. van Hove, CERN-TH-3924, June 1984; *Z. Phys.* **C21**, 193 (1983).
[35] M. Gyulassy, K. Kurki-Suonio, K. Kajantie, L. McLerran, *Nucl. Phys.* **B237**, 477 (1984).
[36] J. Kapusta, Univ. of Minnesota Preprint, Print-84-0134, Jan. 1984.
[37] T. DeGrand, Univ. of Colorado Preprint HU-TFT-84-20, Aug. 1984.
[38] R. Pisarski, *Phys. Lett.* **110B**, 155 (1982); R. Pisarski, F. Wilczek, *Phys. Rev.* **D29**, 338 (1984).
[39] J. D. Bjorken, Fermilab Preprint, Fermilab-Pub-82/59-THY (1982).
[40] T. D. Lee, G. Wick, *Phys. Rev.* **D9**, 2291 (1974).
[41] J. D. Bjorken, L. D. McLerran, *Phys. Rev.* **D20**, 2353 (1979).
[42] A. Kerman, S. A. Chin, *Phys. Rev. Lett.* **43**, 1292 (1979).
[43] E. Witten, *Phys. Rev.* **D30**, 272 (1984).
[44] J. Klosinski, Rivers State University Preprint (Port Harcourt, Nigeria), RSUST-Physics-83/1 (1983).
[45] R. Mackenzie, F. Wilczek, A. Zee, ITP at Santa Barbara Preprint, NSF-ITP-84-99 (1984).
[46] G. L. Shaw, R. Slansky, *Phys. Rev. Lett.* **50**, 1967 (1983).
[47] C. M. G. Lattes, Y. Fujimoto, S. Hasegawa, *Phys. Rep.* **65**, 151 (1980).