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Ultrarelativistic Heavy-ion Collisions and the Quark-gluon Plasma

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Ultrarelativistic heavy-ion collisions and the quark-gluon plasma*

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Abstract. The expected form of nuclear matter under extreme conditions of temperature or baryon density is a quark-gluon plasma. The basic physics of such a plasma and the transition from hadronic to deconfined matter are reviewed. Ultrarelativistic heavy-ion collisions allow one to study properties of nuclear matter under extreme conditions in the laboratory. The formation of plasmas in such collisions and their expected evolution are outlined.

1. Introduction

The basic question I would like to address in this talk is: how does one go about learning the properties of extended matter at extremely high energy densities, an order of magnitude or more beyond that of normal nuclear matter? The energy density of normal nuclear matter is essentially the rest mass density, of order 0.15 $\mathrm{Gev/fm^3}$, large compared with the scale of low-energy spectroscopy, of order $\mathrm{Mev/fm^3}$. How do we expect nuclear matter to act when we raise its energy density to the range of 1-10 $\mathrm{Gev/fm^3}$ say? One of the most interesting possibilities then is that matter will form a new state, the quark-gluon plasma.

To carry out such studies will require colliding heavy nuclei together at energies well above 1 Gev per nucleon in the center-of-mass frame. A program of fixed-target experiments with lighter nuclear projectiles will begin shortly at the CERN SPS at lab energies of 60 and 225 GeV per nucleon, and at the Brookhaven AGS at lab energies of 12-14 GeV per nucleon. These experiments will be reviewed in detail in the following talk by Hans Specht. In addition, the Brookhaven Relativistic Heavy Ion Collider (RHIC) is quite far along on the drawing boards, and will provide the capability of colliding nuclei as heavy as Au on Au at 100 GeV per nucleon c.m. (equivalent to 20 TeV per nucleon lab).[1]

Studying the behavior of matter at high energy densities in the laboratory offers many important opportunities for new physics. Formation of a deconfined quark-gluon plasma will permit one to study quantum chromodynamics (qcd) and quark phenomena over distances very large compared with those available in light hadron systems. In addition, nuclear matter

1. Useful general references on relativistic nucleus-nucleus collisions are the proceedings of the ongoing conferences on quark matter (Ludlam and Wegner 1984, Kajantie 1985, Gyulassy *et al.* 1986).

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under extreme densities is of great interest in astrophysics, a question I will return to later.

The exploration of the behavior of nuclear matter under extreme conditions gives one a rather complementary approach to the study of quark-gluon degrees-of-freedom, and the related question of relativistic effects, in nuclei. High precision electron scattering from nuclei, which will be carried out over the coming years, will provide information on quark properties in individual nuclear states, the effects of quarks in nuclear wave functions. Hadronic probes will provide further detailed information. Ultrarelativistic heavy-ion collisions on the other hand answer a different question: what are the highly excited thermodynamic states, the thermodynamic role of quark and gluon degrees of freedom? As Erich Vogt's comment -- that simply by knowing about Coulomb forces one does not necessarily know the properties of water -- emphasizes, these are rather different aspects of the problem. Heavy-ion collisions will primarily address the problem of the gross statistical behavior of matter, rather than the details of individual wave functions.

Our principal region of interest, in the phase diagram of nuclear matter in the temperature-baryon density plane, Fig. 1, is that at high temperature or baryon density, where one can reach deconfinement of quarks and gluons. In the low temperature-baryon density region the basic degrees-of-freedom are hadronic, those described by Dirk Walecka in the previous talk. Between these two regions may or may not be a sharp phase transition. Later we will come back to the question of how one can learn about the phase diagram from various experiments, but first let us consider the elementary properties of quark matter.

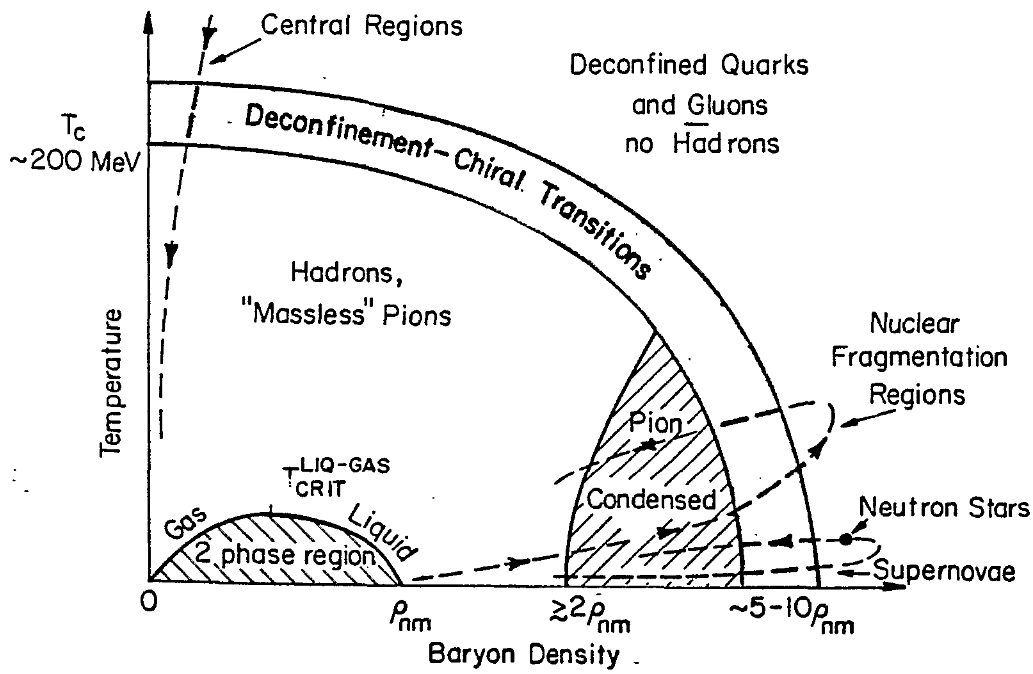

Fig. 1 Phase diagram of nuclear matter in the baryon density, temperature plane showing regions of hadronic and deconfined matter. Normal nuclear matter density ρ_{nm} is 0.16 fm^{-3}.

2. Quark-gluon plasma

As matter is heated or compressed its degrees of freedom change from composite to more fundamental. For example, by heating or compressing a gas of atoms, one eventually forms a plasma in which the nuclei become stripped of the electrons, which go into continuum states forming an electron gas. Similarly, when nuclei are squeezed, as happens in the formation of neutron stars in supernovae where the matter is compressed by gravitational collapse, the matter merges into a continuous fluid of neutrons and protons. Since nucleons themselves are made of quarks, one further expects that a gas of nucleons, when squeezed or heated, turns into a gas of uniform quark matter, composed of quarks, and at a finite temperature, antiquarks and gluons as well, which are no longer confined in individual hadrons but are free to roam over the entire volume of the deconfined region.

Before turning to the properties of quark matter, let us recall a few relevant features of qcd. In quantum electrodynamics, photon exhange produces the basic force between charges, which between static point charges is simply the Coulomb interaction, $\sim e^2/r$. The force between charges of opposite sign is opposite to that between like charges; thus in qed one can form electrically neutral systems, such as positronium or hydrogen, which do not give rise to long-range Coulomb fields. Qcd has a similar structure, in that the forces between quarks arise from exchange of gluons, and the color degree of freedom functions as a three-valued charge, rather than simply + or −, as in qed. Again one can form color singlet or neutral systems that do not give rise to long-range color Coulomb fields. [If one very naively pictures the forces between different colored quarks to be −1/2 that between similar colored quarks, then a nucleon made of three different colored quarks would have no long-range Coulomb interaction with another quark.] To have such a charge scheme requires eight gluons, rather than a single photon, themselves having color and hence coupling directly to themselves producing the rich non-linear structure of qcd. Because qcd allows color neutrality, quark matter in equilibrium in its state of lowest energy, or free energy at finite temperature, will on average have no long-range color Coulomb fields, as in an ordinary electrically neutral plasma.

Qcd is also asymptotically free. In qed an electron gathers around it a polarization cloud of electron-positron pairs in the vacuum which decreases the net charge seen at large distances; the effective charge of an electron at large distances is given by $e^2/\hbar c = 1/137$. At short distances, inside the polarization cloud, the effective charge on the electron grows, diverging at zero distance — one of the troublesome divergences of qed. In qcd a rather different behavior occurs. Because the gluons themselves carry color, they also screen the bare charges, but their net effect is opposite that of quark-antiquark pairs; the result is that as one goes close to a quark, the effective charge does not become infinite, but rather goes to zero — the property of asymptotic freedom. At short distances, corresponding to large momentum scales, interactions become arbitrarily weak [with the effective coupling $\alpha(p) = g^2/4\pi \approx 6\pi/[(33-2N_f)\ln(p/\Lambda)]$ where N_f is the number of quark flavors that are relevant, Λ the qcd scale parameter is of order 100-200 MeV, and p is the momentum scale]. At large distances however quite the opposite happens; if one tries to separate colored particles, the forces become larger and larger, giving rise to confinement.

Imagine then highly compressing or heating nuclear matter, so that many quark-gluon degrees of freedom are excited. To a first approximation, one can at very high densities treat the system as a non-interacting gas of relativistic quarks, antiquarks and gluons, since as in an ordinary plasma, any small region of the matter will, at high densities, be on average color neutral and not produce long-range (color) Coulomb fields, while the residual short distance forces in the region become weak as the interparticle separation becomes small, due to asymptotic freedom. [See Svetitsky 1986 and Polonyi 1986 for a more careful discussion of interaction effects in the high density limit.]

Although for the temperatures and baryon densities of plasmas realistically expected in laboratory collisions, interactions will in fact be important, the non-interacting limit provides a useful first handle on the quark-gluon plasma. While ordinary nuclear matter has 4 helicity states, 2 for spin times 2 for isospin, a quark-gluon plasma has many more internal degrees-of-freedom; the quarks have from 24 to 36 helicity states, composed of 2 spin, 3 color, 2 particle-antiparticle, and 2 to 3 flavor degrees-of-freedom, depending on whether strange (s) quarks are also present in addition to the light up (u) and down (d) quarks; the massless gluons have in addition 16 helicity states (2 spin and 8 color). From a thermodynamic point of view a quark-gluon plasma at a given energy density has a high entropy. Hot free quark matter looks very much like ordinary black-body radiation with energy density $E \sim T^4$, where T is the temperature; in a system with equal number of u, $\bar{u}$, d, and $\bar{d}$, as well as gluons,

$$T \approx 160 \text{ MeV } E^{1/4} \tag{1}$$

and the total density of excitations per fm^3 is

$$n_{exc} \approx 2.25\ E^{3/4} \tag{2}$$

with E measured in GeV/fm^3. Since the qcd phase transition to a quark-gluon plasma is believed to occur at T of order 200 MeV, we see from (1) that the scale of energy densities that must be deposited in collisions to excite a plasma is of order several GeV/fm^3. Because of the slow dependence of T on E, it will not be easy to heat a plasma in a nuclear collision much beyond hundreds of MeV.

In order to neglect interactions, temperatures (or momenta or chemical potentials) should be large compared with the qcd scale parameter Λ. The temperatures produced in collisions are expected, however, to be at most on the order of a few times Λ; one must generally take interactions into account. They are clearly always important near the deconfinement transition. At first one is tempted to use perturbation theory. Writing down all the Feynman diagrams for the thermodyamic free energy one is capable of doing in an afternoon takes one to order $\alpha^2 \ln \alpha$. Taking into account terms of order α^2 is considerably harder. The eventual result, an asymptotic expansion of the free energy in α, is absolutely useless at the coupling strengths of interest. For example, the first correction to the entropy density, calculated by differentiating the free energy with respect to T, is given by $s = s_o[1 - (54/19\pi)\alpha]$, where s_o is the non-interacting entropy density. This first order expression for the entropy turns negative for $\alpha \sim 1$, while coupling constants can become larger at low densities. Now this is impossible — entropies must be positive. What we see is a signal that perturbation theory breaks down early, and is not a satisfactory way to calculate.

The only useful approach so far to calculating effects of interactions is Monte Carlo lattice gauge theory, where by putting the theory of qcd on a lattice one becomes capable of dealing with all strengths of interaction. (See Satz 1985 for a general review.) Lattice gauge theory requires rather large computers to get adequate statistics, but recently, as Cray type supercomputers with substantial megaflop rates have become available to the community, the ability to compute with good statistics on large lattices has improved enormously, and lattice gauge theory is really at the point where it will be able to give quantitatively good information on the properties of quark matter over large ranges of temperature and also baryon density. Calculations with finite baryon density are just in their beginning, and I shall in the limited time available review only the finite temperature, zero baryon density results.

Figure 2 shows early Monte Carlo calculations of the energy density, plotted in units of the ideal non-interacting system (Stefan-Boltzmann) energy, ~ T^4, in a system with just gluons, and no quarks whatsoever -- pure Yang-Mills theory -- in a) for SU(2), with just 2 colors (Engels et al. 1981), and in b) for SU(3) (Celik et al. 1983). The temperature is measured in units of the lattice qcd scale parameter Λ_L. In SU(2), the energy density exhibits a clean second order phase transition. At low temperature the system behaves as a gas of massive glueballs, and is confined, while at the transition temperature, T_c, the system turns smoothly into a gas of deconfined gluons, rather rapidly approaching the high temperature T^4 limit. In SU(3), by contrast, one sees a rather sharp first transition, like the boiling of water, with a large latent heat, of order a few GeV/fm^3. Indeed the Monte Carlo calculations in this case exhibit considerable hysteresis effects, which are a good signal of a first order transition.

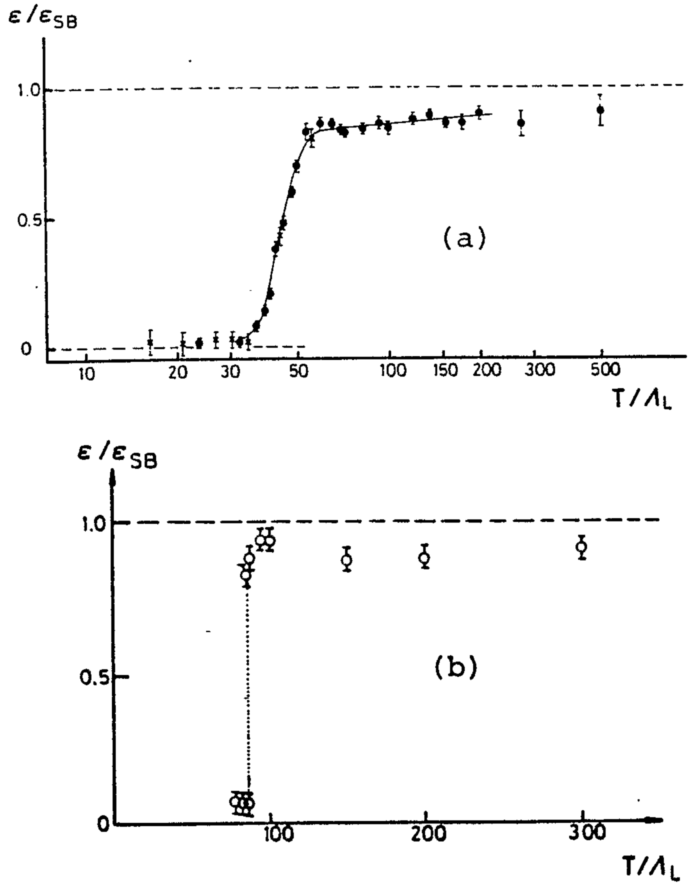

Fig. 2 Energy density of pure gauge theory in units of the ideal gas energy: (a) SU(2), (b) SU(3).

In a pure gluon theory one can measure whether the system is confined or not by adding a massive quark-antiquark pair of test particles to the system, and asking how much energy is required to separate the pair to infinity. If the system is confining, it is impossible to separate them and the energy of separation ε diverges. [Color octet gluons cannot screen the force between color triplet quarks.] Thus the "Wilson line," defined by $W = \exp(-\varepsilon(R\to\infty)/T)$, goes to zero as the separation R goes to infinity, while in the deconfined state it should be non-zero. The Wilson line functions in pure gluon theory as a useful order parameter to distinguish the confined from the deconfined phase. Calculations of W for the pure gluon theory are shown in Fig. 3; in a) for SU(2), W begins to rise above the transition point, corresponding to the onset of deconfinement, while in b) it jumps discontinuously, consistent with the behavior of E.

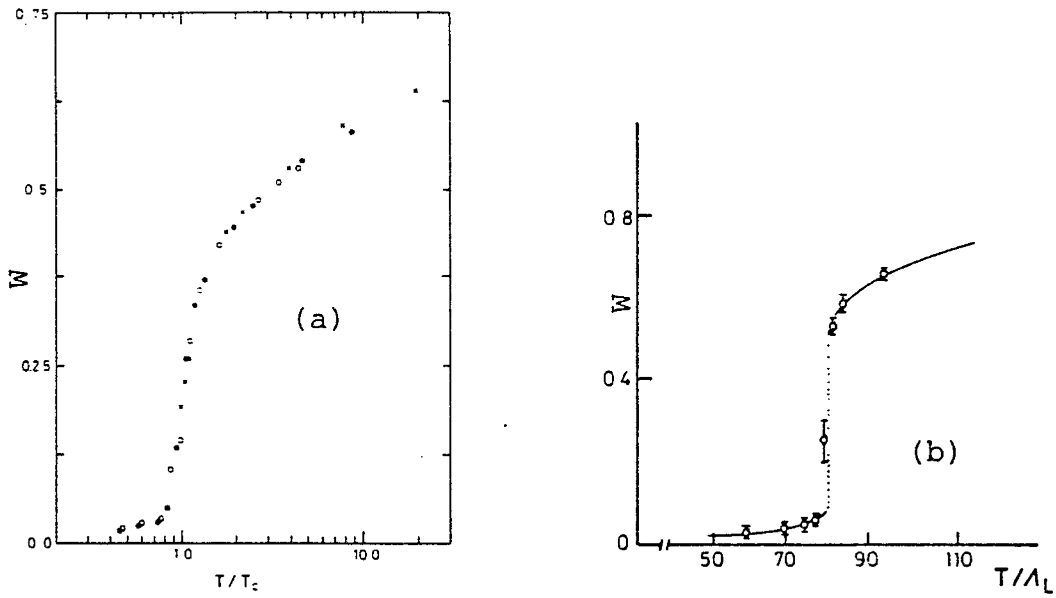

Fig. 3 The Wilson line for (a) SU(2) and (b) SU(3) pure gauge theory.

When one begins to include finite mass quark degrees of freedom, q and $\overline{q}$, the simple test of adding a pair of heavy test quarks $Q\overline{Q}$ runs into trouble, since at sufficient separation, it becomes energetically favorable to create a $q\overline{q}$ pair in the system, which screens out the interaction between the test pair; the q binds to the $\overline{Q}$, and the $\overline{q}$ to the Q, creating effectively a pair of mesons which can be separated to infinity with finite energy. The point is that once one has light quarks in the system there no longer exists a good measure of whether the system is in a confined or deconfined state, and there need not be a sharp transition between the confined and deconfined phases. The transition between the two phases can be smooth, as occurs for example in ionization of a gas as it is heated, where one goes gradually from gas molecules to electrons and nuclei; the two states are qualitatively different and there is a reasonably rapid onset of ionization, but it is not sharp. Alternatively, the transition may be first order, as in the boiling of water.

Qcd matter with light quarks turns out technically to be analogous to a ferromagnet in the presence of an external magnetic field, H, with the quark mass playing a similar role to H^{-1}. Let us consider the case of a ferromagnet in which the transition in zero external field happens to be first order. Then with increasing H the transition becomes weaker and weaker, and finally beyond a critical H_c, the transition becomes smoothed out; in enormous H, the spins are all aligned and there is no transition. In the T-H plane, one finds a line of first order transition points, terminating in a critical point at H_c. Similarly a quark-gluon plasma, at large quark mass, has a well-defined transition, which weakens as the mass decreases; it is not immediately apparent whether for realistic quark masses, analogous to large H, there is a sharp transition or not. One problem is that to take light mass quarks into account accurately in lattice gauge theory requires quite extensive computing, and only very recently has one begun to achieve a reasonable understanding of the nature of the transition (at zero baryon density).

One further ingredient which must be taken into account is chiral symmetry, the nearly exact SU(2)xSU(2) symmetry of the strong interactions in low nuclear physics generated by the conserved vector current

$$\vec{V} = \int d^3r\ \overline{\psi}(r)\gamma^o\vec{\tau}\psi(r)$$

together with the partially conserved axial vector current (PCAC)

$$\vec{A} = \int d^3r\ \overline{\psi}(r)\gamma^o\gamma_5\vec{\tau}\psi(r)\ .$$

Because the axial current is not precisely conserved, chiral symmetry is not exact; the level of violation is measured by the smallness of the pion mass: $(m_\pi/m_n)^2 \sim 1/50$. From the point of view of the underlying quark structure, chiral symmetry is exact only for zero mass u and d quarks; the violation of chiral symmetry is a reflection of the fact that the light quark are not precisely massless, but have masses on the order of 10 MeV.

Symmetries can be realized in a physical situation in two ways: the Wigner mode in which the states can be classified according to the representations of the rotation group, as in atoms and nuclei, and the Goldstone mode, in which the equilibrium state picks out a given direction in the group space, analogous to the situation in a ferromagnet (or more precisely an anti-ferromagnet), where the magnetized state chooses a special spatial direction for the magnetization, breaking the overall rotational symmetry of the state. In nuclear physics chiral symmetry is in fact spontaneously broken. The long wavelength small oscillations of the spins in an aligned ferromagnet are low-lying modes, the spin waves; in the case of broken chiral symmetry, the analogous low-lying excitations, or Goldstone bosons -- the oscillations of the spontaneously selected direction -- are the physical pions.

The implications for qcd are the following. At low temperatures and baryon densities chiral symmetry is spontaneously broken. On the other hand, if at the high momentum scales of very high temperatures or densities the system becomes asymptotically free, then one expects chiral symmetry to be fully restored. Between these two limits a chiral symmetry restoring phase transition should occur. Associated with this transition is a well-defined order parameter $\langle\overline{\psi}\psi\rangle$ (where ψ is the quark field) which is non-zero in the spontaneously broken phase and zero where chiral symmetry is fully restored. Finite quark mass can wash out the chiral transition; however as m becomes sufficiently small one expects the appearance again of a sharp transition associated with the chiral symmetry.

Figure 4 shows several recent calculations of the Illinois Monte Carlo group (Kogut and Sinclair 1986, and earlier references therein) in SU(3) with finite quark mass m, on a 6x10x10x10 lattice. The horizontal scale is the effective coupling $\beta = 6/g^2$, which increases monotonically with temperature. The mass is given in terms of the transition temperature T_c (generally of order twice the qcd scale parameter Λ) by 6 times the values in the figures (there in units of a the lattice spacing). Fig. 4a, the Wilson line for a system with relatively heavy quarks -- over an order of magnitude more massive than the u and d quarks -- nicely illustrates how matter with heavy quarks has a first order phase transition. The Wilson line shows a fairly sharp onset of deconfinement; hysteresis in the transition provides good evidence that it is actually first order. Figure 4b, for intermediate mass $ma = 0.05$, shows a smooth onset of the deconfined phase; the behavior is more like a second order or smoothed-out phase transition. Also plotted is the chiral order parameter $\langle\bar{\psi}\psi\rangle$, which goes rapidly from a finite to a small value (not exactly to zero, since the quark mass is finite in these calculations) as deconfinement sets in.

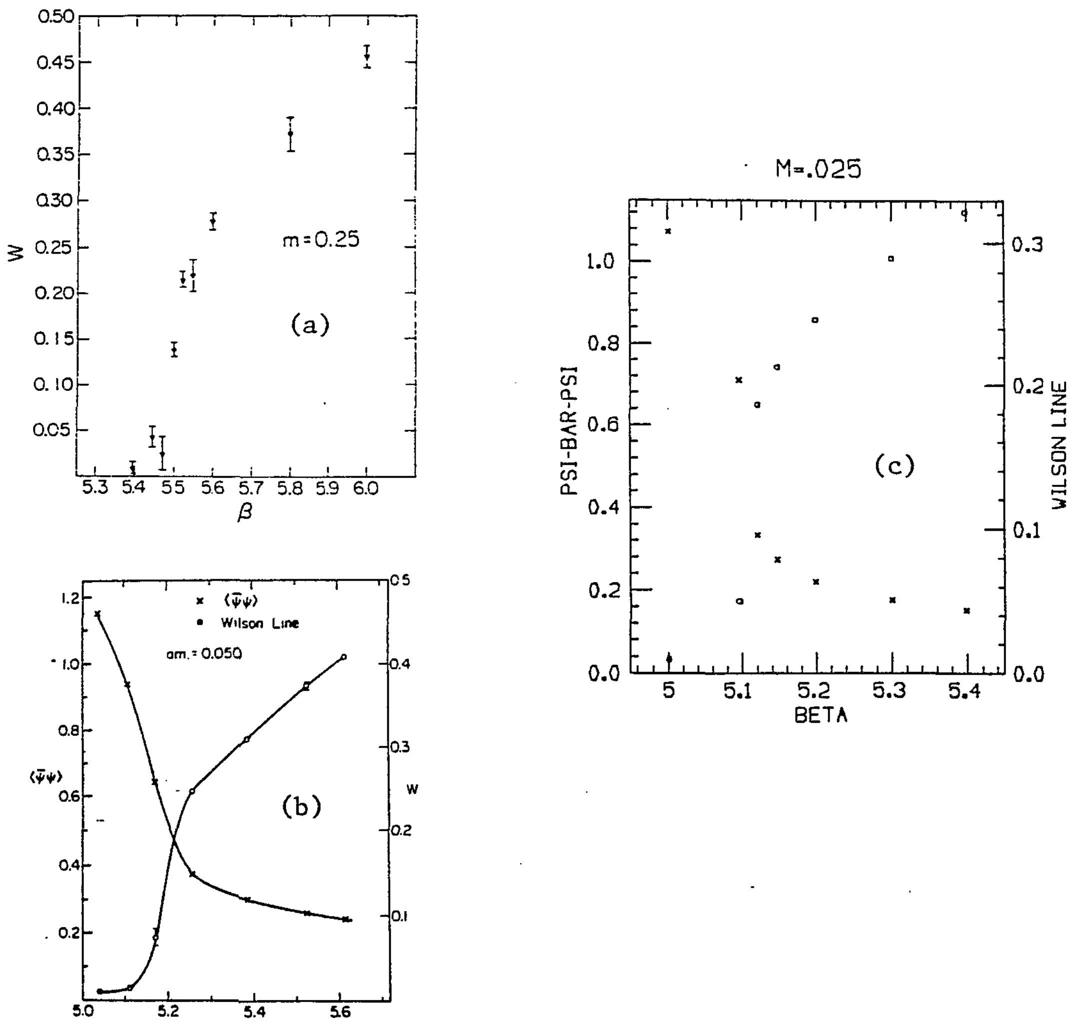

Fig. 4 Wilson line in SU(3) with a) relatively heavy quarks, b) intermediate mass; c) light mass; b) and c) show the chiral behavior as well.

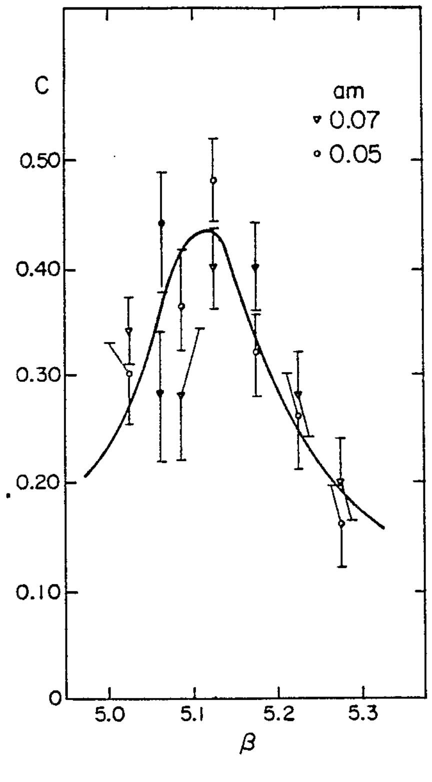

Fig. 5 Specific heat with intermediate mass quarks.

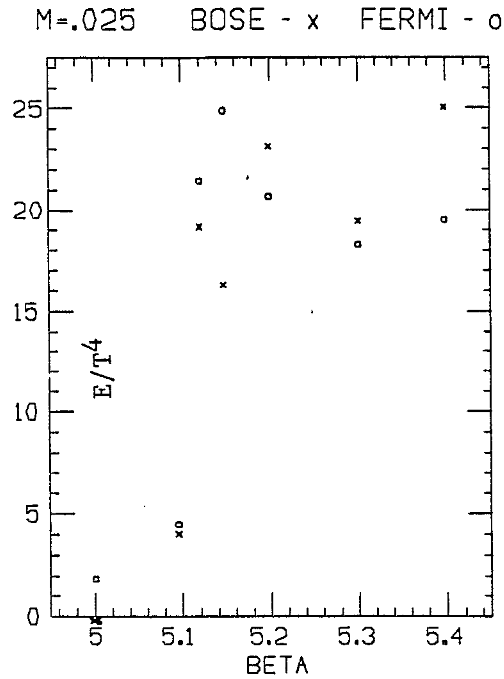

Fig. 6 Energy density in units of T^4 for light quarks.

Another measure of the sharpness of the transition is the specific heat, which indicates the temperature range over which energy must be put into the system to go from one phase to the other; Fig. 5 (Kogut 1986) shows, for the intermediate mass cases ma = 0.05 and 0.07, the specific heat versus β (or effectively the temperature) in the neighborhood of the transition. We again see not a sharp transition in this case, but rather a large but smooth bump; the total energy under the curve is of order a few GeV/fm^3.

However for even lower mass, in Fig. 4c (Kogut and Sinclair unpublished) -- ma = 0.025, $m/T_c = 0.15$, closer to the realistic case -- we see evidence that the restoration of chiral symmetry, and deconfinement, is now via a sharp transition. Figure 6 (Kogut and Sinclair, to be published) shows the energy density divided by T^4, versus temperature for this light mass case. Above the deconfinement transition, the energy goes rapidly to the Stefan-Boltzmann limit (19.1 for fermion quarks with 4 flavors, as considered here).

To summarize, present lattice gauge theory calculations with finite mass quarks indicate that at large m the transition between the confined and deconfined phases is first order, while with decreasing m one enters an intermediate regime where the transition is somewhat washed out; at still lower quark masses, when the chiral behavior is accurately taken into account the transition sharpens again, driven by restoration of chiral symmetry, and is likely (weakly) first order.

3. Exploring the phase diagram

The physics of matter at high temperature and density can be explored in astrophysical as well as laboratory situations. Let us first briefly discuss the applications to supernovae, neutron stars and the early universe. In a (Type II) supernova explosion, a massive star which has burnt out its fuel at the end of its evolution can no longer support itself against gravitational collapse, and begins to implode. In the infall the matter is crushed to very high densities, several times that of normal nuclear matter. The core of the star bounces back, as shown by the trajectory in the phase diagram; the matter in the core may or may not cross into the deconfined region. The energy output in a supernovae, which should depend on the strength of this bounce, may provide a handle on the nature of the matter in the interior reached in the collapse.

In neutron stars the properties of matter under extreme conditions play a particularly crucial role. One quickly learns in trying to construct models of neutron stars how little is known about the properties of matter at densities beyond $\sim 2\rho_{nm}$. The lack of such knowledge is reflected, for example, in uncertainty in the maximum mass of neutron stars, an important quantity in trying to identify black holes unambiguously. Measured masses of neutron stars are ~ 1.4 solar masses, with radii calculated to be ~ 10 km. Neutron stars are giant nuclei, with $A \sim 10^{57}$. Typical temperatures are very low, less than one MeV. The central conditions in a neutron star are indicated on the phase diagram. The matter in the interior may possibly be deconfined.

One can in fact study the interior of neutron stars by observing their cooling. In their early years ($<10^5$y), cooling is governed primarily by neutrino emission. For phase space reasons, cooling via the nucleonic URCA process, $n \rightarrow p + e + \bar{\nu}$, and $e + p \rightarrow n + \nu$, is considerably slower than it would be via the corresponding process with light mass deconfined quarks, $d \rightarrow u + e + \bar{\nu}$, and $u + e \rightarrow d + \nu$. Combining knowledge of the ages of astrophysical objects containing neutron stars with measurements of neutron star surface temperatures, taken with X-ray telescopes, gives one a measure of how rapidly neutron stars cool. Particularly rapid cooling would provide evidence for unusual states of matter in the interior. Present observations, which generally provide an upper bound on surface temperatures, are so far consistent with the interiors being normal nuclear matter, although future satellite observations should sharpen these bounds and provide a more definitive answer to the nature of the matter in the interiors of neutron stars.

In the first microseconds of the early universe, the temperature falls as

$$T \approx 0.5\ (t_{\mathrm{seconds}})^{-1/2}\ \mathrm{MeV}\ ,$$

so that prior to ~ 6 microseconds after the big bang, when the temperatures are > hundreds of MeV, matter is in the form of a quark-gluon plasma. The matter of the early universe has a relatively small net baryon density, of order 1 part in 10^9 (which appears as the present photon/baryon ratio). As the universe expands it cools and matter hadronizes, following a downward trajectory practically along the vertical axis of the phase diagram. Matter emerging from the transition is primarily in the form of pions, with a slight baryon excess. Possible astrophysical consequences of the transition in the early universe from deconfined plasma to hadrons will be touched upon by Dave Schramm in his talk.

4. Ultrarelativistic heavy-ion collisions

Coming back down to earth, let us consider how one can study the phase diagram in the laboratory via very energetic collisions of heavy nuclei. Imagine colliding two nuclei together at energies from 1 to 100 GeV per nucleon energy in the center-of-mass. At low energies, the regime that will be studied in the forthcoming SPS and AGS experiments, one can picture the two Lorentz-contracted colliding nuclei as nearly stopping each other, with reasonable probability of forming, to a crude first approximation, a fireball. [In reality, parts of the nuclei will generally pass through the collision rather than remain in a fireball.] Such collisions may achieve energy densities of order a few Gev/fm^3 and baryon densities several times ρ_{nm}; the matter may indeed cross into the deconfined region, as shown by the curve in Fig. 1 labelled "fragmentation regions," and then expand out. The high density matter produced in such collisions will be relatively baryon rich.

As the beam energy is increased, the time τ_0 it takes for excitations to form as the nuclei collide, measured in the center-of-mass, becomes, as a consequence of Lorentz time dilation, effectively longer than the time it takes for the nuclei to pass through each other -- the phenomenon of nuclear transparency. In this case the nuclei pass through each other, become highly excited internally, and at the same time, leave the vacuum between them in a highly excited state containing quarks, antiquarks and gluons as illustrated in Fig. 7. Nuclear transparency is very important in the ultrarelativistic regime, above ~ 10 GeV per nucleon (c.m.). The nuclear fragmentation regions, which recede from each other at the speed of light, contain essentially all the baryons of the original nuclei; the central region, to a first approximation, has no baryon excess, and resembles the hot vacuum of the early universe. [In fact, the baryons of the colliding nuclei will spread somewhat into the central region; predicting how much is an important problem on which the forthcoming heavy-ion experiments at CERN and Brookhaven are expected to shed light. However, the energy per baryon will in general be very high.]

In ultrahigh energy collisions, one expects a useful strong correlation between the spatial structure of the collision region, and the final rapidities of the detected particles emerging from the collision.

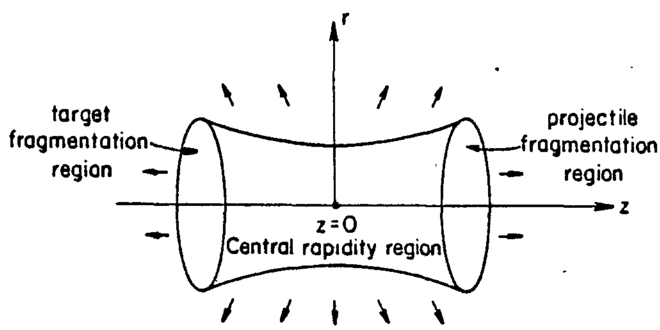

Fig. 7 Nuclear fragmentation and central rapidity regions in an ultrarelativistic central heavy-ion collision.

[Relativistic collisions are conveniently described in terms of the rapidity y of the beams and collision products, defined by $y = (1/2)\ln[(E+p_z)/(E-p_z)]$, where E is the particle energy, and p_z its momentum along the beam axis. For motion purely along the z axis, the velocity is given by $v/c = \tanh y$. Rapidities, unlike ordinary relativistic velocities, have the nice property of being additive; under a Lorentz transformation along z by velocity $u = c \tanh y_u$, rapidities transform by $y \to y + y_u$.] To see how this correlation works, think of the fragmentation regions after the collision as travelling at c away from each other, with the central region being uniformly stretched out in between, so that its velocity increases linearly with distance, from −c at the left fragmentation region to +c at the right fragmentation region (Fig. 7). Thus, particles observed at large positive rapidities come primarily from the right fragmentation region, those at large negative rapidities from the left fragmentation region, while the intermediate rapidity particles arise from the central regions. The intrinsic motion of the particles with respect to the local average motion will, of course, blur this correspondence somewhat, on the order of one unit of rapidity.

The total spread in rapidity in a collision is given by $\Delta y = 2 \ln(2E/m_n)$ where E is the beam energy per nucleon in the center-of-mass. Thus ultrarelativistic collisions provide sufficient total rapidity spread -- for example 10.6 units at 100 GeV on 100 GeV -- that this correlation enables one to sort out the different collision regions from the rapidities of their final state products; this ability to distinguish different regions experimentally is one of the principal reasons for going to high energies. The baryons will appear predominantly at large rapidities, while in the central rapidity region, which has little excess of baryons over anti-baryons, one will see primarily mesons. Indeed, such a structure emerges in proton-proton scattering, as carried out at the CERN ISR, and in $\bar{p}p$ collisions at the $S\bar{p}pS$. The charged particle multiplicity vs. rapidity for 30 GeV on 30 GeV pp collisions is shown schematically in Fig. 8; the meson spectrum (unshaded) is spread out over the central region, while the net baryon density (shaded) is peaked near the rapidities of the two incident colliding beams, with a width of order 2 units in rapidity.

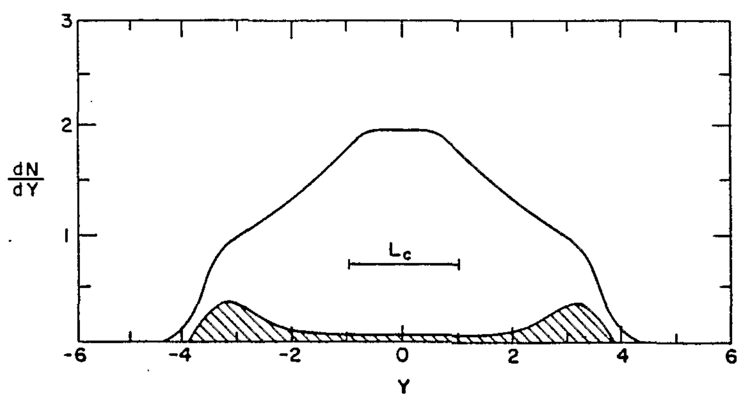

Fig. 8 Charged particle multiplicity in pp collisions at 30 GeV on 30 GeV.

The scale of energy densities expected in ultrarelativistic collisions can be estimated from the observation that pp or $\bar{p}p$ collisions in this energy range produce ~ 2-3 charged particles, predominantly pions, per unit of rapidity, with typical transverse energy, ~ 400 MeV. Adding in neutral pions as well, we find an energy density of ~ 2 GeV per unit of rapidity. A very conservative extrapolation to a nucleus-nucleus collision is to multiply this produced energy by a factor ~ (1-2) A, i.e., assume that each nucleon of one of the nuclei makes only one or two collisions going through the other nucleus. The resulting estimate of the energy per unit rapidity in an AA collision is > (2-4)A GeV, and the corresponding energy density is

$$E \gtrsim 0.4A^{1/3}/t \text{ GeV/fm}^3$$

where t is the time in fm/c. At a time of 1 fm/c the energy density is at least of order 2.5 GeV/fm^3 in the average central collision.

The actual energy density can be much larger, as can be seen from a similar argument based on extrapolating from particle multiplicities (Von Gersdorff et al. 1986). If we assume that entropy is conserved in the evolution of the collision, then at early times, $tT^3 \sim (dN/dy)/\pi R^2$, where T is the temperature, R is the nuclear radius, and dN/dy is the final multiplicity density in rapidity. Clearly the earlier the time, the greater the temperature and hence energy density. The uncertainty principle implies that the initial formation time, τ_o, when one can first begin to describe the system in terms of well-defined interacting excitations (quarks, antiquarks and gluons), and the initial temperature, T_o, at that time, obey $T_o\tau_o \gtrsim 1$. Let us assume that the mean multiplicity in an AA collision is of order A times that in a pp collision. Combining with the uncertainty principle relation we deduce that the formation time is $\sim A^{-1/6}$ in fm/c, that the initial temperature is ~ 200 $A^{1/6}$ MeV, and most importantly, the initial energy density one expects to achieve, proportional to $T_o{}^4$, is $E \sim A^{2/3}$ in GeV/fm^3, which is of order 30 GeV/fm^3 for Au or U collision partners. Furthermore, in rare events, on which one can certainly trigger, one expects even larger energy depositions. Thus we have good reason to believe that the energy densities in collisions will be sufficiently high in many events to form interesting states of matter.

A further feature of pp and $\bar{p}p$ scattering observed at CERN collider energies is that the multiplicity distributions in the central regions are roughly flat, as can be seen in Fig. 8. Since changing rapidity is equivalent to making a Lorentz transformation, the lack of change under a shift of the horizontal rapidity scale implies that conditions in the central region are approximately Lorentz invariant. The assumption of Lorentz invariance in the central region provides a very simple first picture of the evolution of the central region, as is illustrated in the space-time diagram, Fig. 9, which shows a slice of the collision along the central axis, z, versus time t. At negative times the projectile and target approach each other along the light cone and collide. The first event in the collision is production of excitations, after a finite formation time, which stream out from the collision point, the origin. In heavy nucleus-nucleus collisions, unlike in pp collisions, sufficiently large numbers of excitations are made, and the system sizes are correspondingly large, that after a further finite time the excitations come into local thermodynamic equilibrium; this means that the system enters a regime where it can be described by hydrodynamics. Two interesting phenomena occur now. One is that the matter passes through the

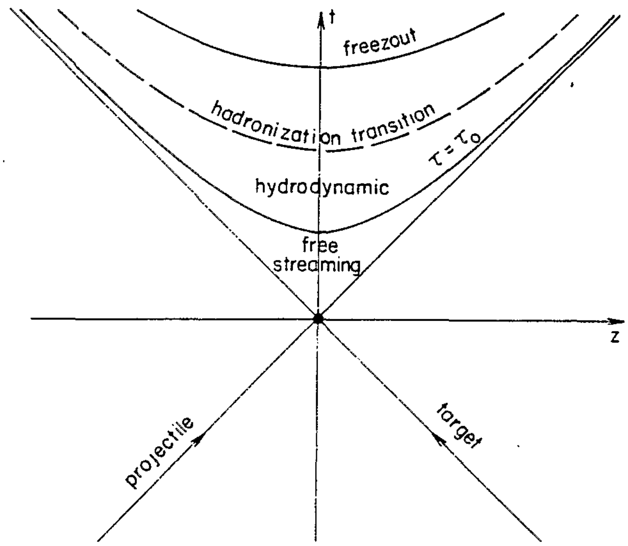

Fig. 9 Space-time picture of the evolution of a central collision.

hadronization transition, as shown in the phase diagram by the curve labelled "central regions," and emerges from the transition in the form of hadrons. Eventually the system expands sufficiently that the interactions among the hadrons cease -- "freezeout" -- and the system becomes a collection of freely streaming particles, which are eventually detected. The lines separating the regions in the space-time diagram are essentially hyperbolae, since as a consequence of the approximate Lorentz invariance in the central region, one expects the same conditions at all z in the central region at the same proper time, $(t^2 - z^2)^{1/2}$ (corresponding to local velocity z/t along the central axis).

One very important point that this diagram illustrates is that the matter undergoes considerable processing from the initial quark-gluon plasma phase to the finally observed hadronic products; tracing back from the observations to the properties of the plasma will clearly be a challenging problem. In the collision, the two fragmentation regions move away from each other longitudinally, and the central region undergoes a longitudinal stretching and cooling, as illustrated in Fig. 9. In addition, the system begins to undergo transverse expansion, which initially occurs hydrodynamically with a rarefaction wave propagating inward from the outer edge at the speed of sound.

To summarize, the basic picture of ultrarelativistic collisions, while clearly a large extrapolation from measured pp and $\overline{p}p$ collisions, is likely rather well understood on the whole. The situation is reminiscent of the recent discovery in New Mexico of some 8 fragments of data -- which turned out to be dinosaur tail bones -- in a cliff. With further digging

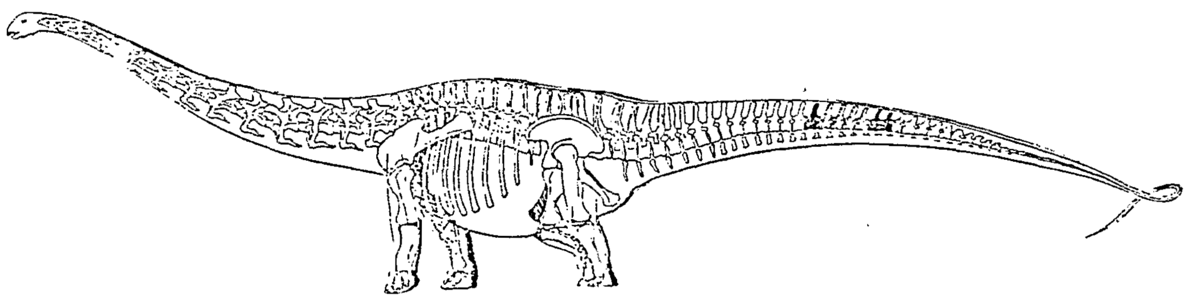

Fig.10

several leg bones were discovered, from which a very beautiful extrapolation was made in Fig. 10 -- the world's largest dinosaur. [The bones that were found are shaded.] The extrapolation is really remarkable, even to the smile on the face and the curl of the tail. These details may be wrong: the neck may go straight up and the dinosaur may face the other way and be frowning, but the basic outlines of the picture, like those of ultrarelativistic collisions, are reasonable. In ultrarelativistic collisions, we not only have the possibility of extracting the physics of the states of matter produced in collisions from the forthcoming data (albeit a difficult task), but we have the possibility of discovering major surprises, which together make the study of ultrarelativistic collisions very exciting.

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