DECONFINED PHASES OF STRONGLY INTERACTING MATTER

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ABSTRACT

This talk surveys our current understanding of deconfined phases of strongly interacting matter at high energy density -- quark matter, or quark-gluon plasma -- as is likely to be produced in ultra-relativistic heavy ion collisions. Following a review of aspects of the normal and perturbative vacua in quantum chromodynamics, elementary properties of deconfined matter and ways in which this state can be explored in both laboratory and astrophysical settings are described. Finally, the possible existence of an intermediate deconfined massive quark phase, in which chiral symmetry remains spontaneously broken, is discussed.

1. Introduction

Ultra-relativistic heavy ion collisions present the enticing possibility of allowing us to study extended matter in the laboratory under extremes of energy density reachable otherwise only in the hot early universe, or in certain astrophysical situations of strong gravitational collapse. In particular, such collisions can possibly create regions of matter so energetic that the normal forces that confine quarks and gluons into individual hadrons are

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overcome, and the matter in the interaction region is in the form of an extended quark-gluon plasma. In this talk I would like to give an overview of the possible states and properties of deconfined quark matter that one may hope to produce in heavy ion collisions. [The confinement of this talk itself to the present volume limits the detail with which many of the topics here are discussed, and the reader is referred to earlier reviews[1-3] for more leisurely treatments.] In order to understand the deconfined phases, let us set the stage with a brief review of the several states of the vacuum in quantum chromodynamics (qcd).

2. States of the Vacuum

On the short distance scales within hadrons, asymptotic freedom implies that the effects of qcd are small and can be treated perturbatively; on the other hand, on the large distance scales outside hadrons, interactions are very strong and drastically affect the properties of the vacuum. In the absence of a complete solution of qcd, one can describe confinement and possible deconfinement of quarks to a first approximation in terms of a picture of the vacuum having two possible phases. The first, the normal vacuum outside hadrons, is that in the absence of physical quarks and color fields. Quarks and gluons modify the vacuum in their neighborhood, transforming it into a second, higher energy state, the perturbative vacuum, the form of the vacuum inside hadrons. The crucial difference between these two states is that the normal vacuum excludes physical quark and gluon fields, while they can propagate freely throughout the perturbative vacuum. In terms of quark masses, one would say that in the normal vacuum the mass of an isolated quark is infinite (provided confinement is exact), while in the perturbative vacuum the quarks have the current mass values, $\lesssim$ 10 MeV for u and d quarks (in order for the chiral symmetry of strong interactions to be so nearly exact) and ~300 MeV for the strange quark. [By contrast, the constituent or effective masses of light quarks in hadrons are ~300 MeV.]

The picture of the vacuum having two states is the basis of the MIT and related bag models of hadrons. The higher energy density B of the perturbative vacuum, compared with the normal vacuum, is essentially the "bag constant" of these models (although its value may

differ from the phenomenological value, $B \approx 56$ MeV/fm^3, for example, deduced from fitting hadron masses in the MIT model). In a bag model, the quarks composing a hadron are assumed to transform a finite region of space to which the quarks are confined, the bag, from normal to perturbative vacuum. The energy density B then acts as an effctive pressure of the normal vacuum on the bag.

A very appealing characterization of the confining property of the normal vacuum, the exclusion of quark and gluon fields, is in terms of the color dielectric constant ε of the two vacuum states.[4,5] The dielectric constant is defined, as in Maxwell, by the relation

$$\underset{\sim}{D} = \varepsilon \underset{\sim}{E} \tag{1}$$

between the color D field, the field produced by the source charges, and the color E field, the total color field including that produced by induced charges. Then one can characterize the perturbative vacuum by $\varepsilon = 1$, so $\underset{\sim}{D} = \underset{\sim}{E}$, and the normal vacuum by $\varepsilon = 0$. The vanishing of the dielectric constant produces confinement of color D fields, since the total energy density of a physical field $\propto \underset{\sim}{D}\cdot\underset{\sim}{E} = D^2/\varepsilon$ can be finite only if $D = 0$ in the normal vacuum.

As emphasized by Lee,[5] the vanishing of ε in the normal vacuum leads to confinement of quarks, as well as color fields. One may see this by computing the electrostatic self-energy of a single charge of strength g, assuming that in the normal vacuum $\varepsilon < 1$, that the charge converts its immediate neighborhood (a sphere of radius R; see Fig. 1) to normal vacuum with $\varepsilon = 1$, and that it costs energy B per unit volume to have $\varepsilon = 1$. Suppressing the SU(3) color algebra we have

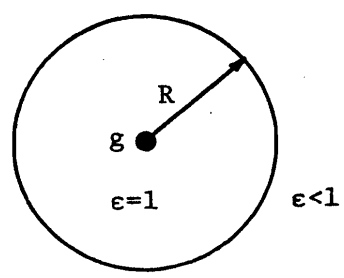

Fig. 1

$$\underset{\sim}{D} = g\hat{\underset{\sim}{r}}/4\pi r^2, \qquad \underset{\sim}{E} = g\hat{\underset{\sim}{r}}/4\pi r^2\varepsilon(r) \ , \tag{2}$$

since the normal component of D is continuous across the boundary at

R. Then the energy of the configuration is given by

$$E = \frac{4\pi R^3}{3} B + \frac{1}{2} \int_a^\infty d^3r \, \underset{\sim}{D} \cdot \underset{\sim}{E}$$

$$= \frac{4\pi R^3}{3} B + \frac{g^2}{8\pi} \left[\frac{\varepsilon^{-1}-1}{R} + \frac{1}{a} \right] , \qquad (3)$$

where a is a small cutoff. If $\varepsilon > 1$, the energy is minimized by having R be as small as possible, while if $\varepsilon < 1$ the energy has a minimum at a finite $R \propto (\varepsilon^{-1}-1)^{1/4}$; hence at the minimum $E \propto (\varepsilon^{-1}-1)^{3/4}$, which diverges as $\varepsilon \to 0$. In other words, the total energy, or mass, of a single charge becomes infinite as $\varepsilon \to 0$, so that a single charge cannot exist in isolation but rather must be confined by forming charge neutral combinations with other charges. The infinite energy arises because in order to satisfy Gauss' law, the field of a single charge must have a long ranged component. In contrast, a neutral pair of charges has only a short ranged field, and a finite energy as $\varepsilon \to 0$. Thus, vanishing of ε in the normal vacuum provides a simple mechanism for both field and quark confinement.

The dielectric behavior of the normal vacuum is analogous to the perfect diamagnetism of a superconductor, which when placed in a magnetic field expels the magnetic flux from its interior -- the Meissner effect. This property can be described by a vanishing of the permeability μ, defined by $\underset{\sim}{B} = \mu \underset{\sim}{H}$, inside the superconductor. Outside a superconducting cylinder placed in a diamagnetic field with the cylinder axis parallel to the field, one has $\underset{\sim}{B} = \underset{\sim}{H}$, while $\underset{\sim}{B} = 0$ (although $\underset{\sim}{H}$ is non-vanishing inside the superconductor, since the tangential component of $\underset{\sim}{H}$, and hence $\underset{\sim}{H}$ in this geometry, is continuous across the surface). In qcd the D field is expelled from the normal vacuum -- an "electric" Meissner effect. If in a (type I) superconductor the external field exceeds a critical value H_c the material becomes normal; analogously one expects a critical field D_c in qcd in the sense that at an interface between normal and perturbative vacuum the mean D field cannot exceed D_c, since then the stress $D^2/2$ of the field would exceed the inward pressure of the normal on the perturbative vacuum, turning the normal vacuum into the perturbative state.

While the confining properties of the states of the vacuum can be described in terms of their color dielectric properties arising from the underlying structure of the color fields in the vacuum, measure-

ment of the dielectric properties by the usual procedure of determining the force between test charges is, however, elusive, since if one puts in a pair of test quarks in the confining vacuum, at large separation they will, because of the possibility of $q\bar{q}$ pair creation, have their charges perfectly screened and will experience no long distance force. One may ask whether there is a more fundamental order parameter distinguishing the two phases. For example in a (BCS) superconductor, rather than the local magnetic permeability, μ, the fundamental order parameter is the off-diagonal amplitude $\langle\psi_\uparrow(\underset{\sim}{r})\psi_\downarrow(\underset{\sim}{r})\rangle$ for removing a pair of electrons, one with spin up and one with spin down, from the ground state, and returning to the ground state; in the normal state this order parameter vanishes. [The permeability μ, and hence a magnetic field, falls to zero only over a distance of a penetration depth in from the surface of the superconductor, although within the penetration depth the material is superconducting.]

A possibly more fundamental characterization of the two states of the vacuum is in terms of the way they realize the SU(2)×SU(2) chiral symmetry corresponding to conservation of the vector current and (partial) conservation of the axial current. In the normal vacuum chiral symmetry is, as we know, spontaneously broken, and realized in the Goldstone mode; the corresponding Goldstone boson is the pion, whose small mass ($m_\pi^2/m_n^2 = 1/50$, where m_n is the nucleon mass) is a measure of the accuracy of the chiral symmetry of strong interactions. This realization of chiral symmetry is analogous to the way rotational symmetry is realized in a ferromagnet, where the magnetization vector picks out a particular direction in space. On the other hand, asymptotic freedom strongly suggests that chiral symmetry is fully restored, i.e., realized in the Wigner mode, in the perturbative vacuum; the light quarks, u and d, are essentially massless in this state. The two phases would then be characterized by an order parameter $\sim \langle\bar{q}q\rangle$ (where q is the light quark field operator), which vanishes in the perturbative vacuum, but is non-zero in the normal vacuum. There is certain evidence, which we will discuss later, that in the transition, with increasing temperature or baryon density, from the normal to perturbative vacuum, the deconfinement transition somewhat precedes the chiral transition; generally however, the connection between the two characterizations of the states of the vacuum in terms

of their dielectric and chiral properties remains an important and difficult problem. [For further discussion of chiral symmetry and bag models of hadrons see Ref. 6.]

3. Deconfinement at High Temperature or Baryon Density

When the density of quarks plus anti-quarks in a system is low, the quarks are confined in individual hadrons, surrounded by normal vacuum. However, as the density is raised, either by increasing the baryon density, or producing hadrons thermally with increasing temperature, the hadrons begin to overlap and the matter is expected eventually to undergo a transition to a state -- quark matter, or quark-gluon plasma -- in which the quarks and gluons are no longer locally confined, but are free to roam over the entire system. If one imagines hadrons as surrounded by little islands of perturbative vacuum, as in bag models, then at sufficiently high density, the in-between regions of normal vacuum are squeezed out, and the space becomes filled with perturbative vacuum.

To a first approximation we expect the phase diagram of matter in the baryon density n_b, temperature T, plane to have the structure shown in Fig. 2. At low n_b and T matter is composed of confined hadrons, but as the energy density is raised, with increasing T or n_b, or both, matter becomes deconfined. For example, if one could drop a proton into the hot vacuum with $T > T_{dec}$, the $n_b = 0$ deconfinement temperature, it would evaporate into three quarks. The temperature T_{dec} has been estimated, by a variety of ways discussed below, to be ~ 200 MeV, and the critical density n_{dec} for deconfinement at T = 0 of the order of a few to more than 10 times nuclear matter density (n_o = 0.16 fm^{-3}).[2] However the order of the deconfinement transition is uncertain. If the chiral occurs after the deconfinement transition, as Monte Carlo calculations indicate,[7,8] then the lack of a satisfactory order parameter to describe confinement suggests that the deconfinement transition may smoothed out or perhaps be first order.

The deconfined phase of matter can be explored in a variety of ways, as indicated on the phase diagram. In addition to man-made production of quark matter in heavy ion collisions, nature provides several laboratories in which the presence of quark matter could be detected, including, cosmic rays, supernovae, neutron stars, and the

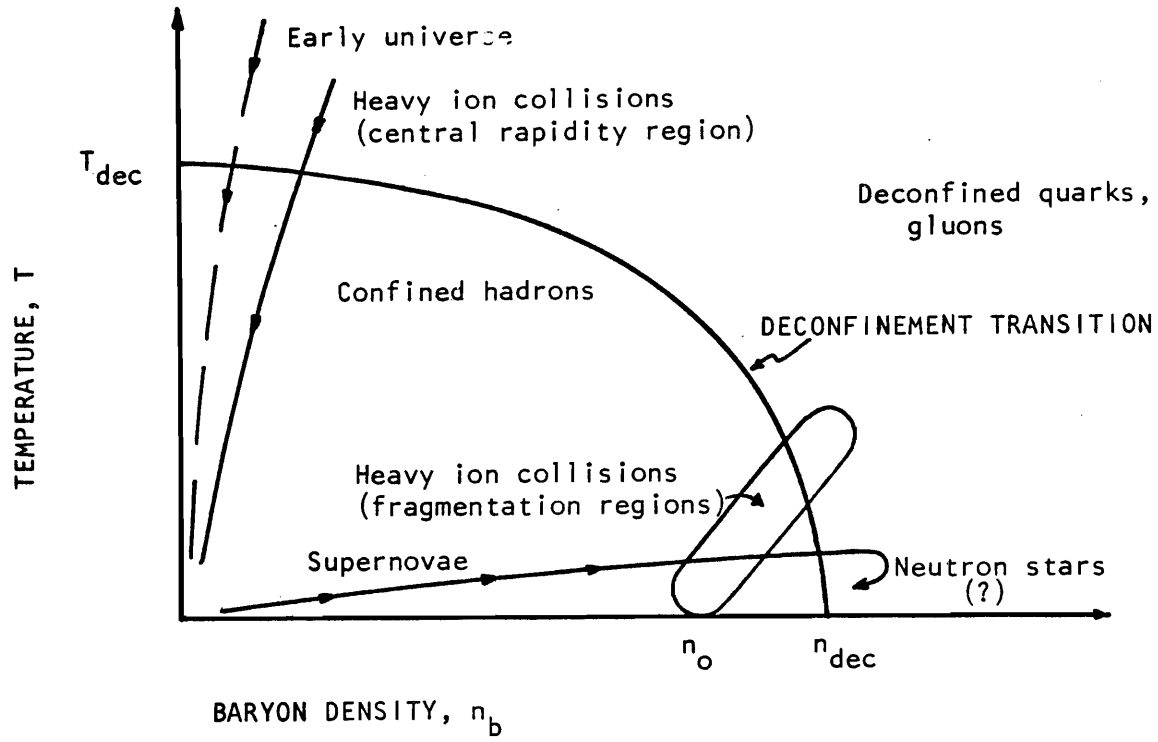

Fig. 2. Phase diagram of matter in the temperature, baryon density plane, with indication of the regimes that might be reached in different physical situations.

early universe. Among cosmic ray phenomena that could involve large droplets of quark matter[9] are the spectacular Centauro events.[10] In supernovae, the explosions of more massive stars at the end of their evolution, the central region is highly compressed on millisecond time scales, through gravitational collapse and forms, it is believed, a neutron star. [For an overview see Ref. 11.] The central matter, which is in thermal equilibrium, follows a trajectory as shown, with the temperature remaining less than ~MeV, rather cold on a scale of nuclear collisions, and the baryon density rising to at least several times n_o, possibly crossing the threshold for deconfinement. Supernovae can, in principle at least, provide information on the nature of the dense matter reached in the collapse, as reflected through the total energy released, for example.

Neutron stars, which are relatively cold ($T \lesssim 1$ MeV), may have central baryon densities above the deconfinement density. The existence of quark matter cores in neutron stars could be detected from observations of their cooling (reviewed in Ref. 1). In their first 10^6 years or so neutron stars cool primarily by neutrino emission

processes from their interiors. The essential point is that, because of phase space considerations, degenerate massless quark matter cools much more rapidly via the neutrino URCA process,[12)]

$$d \to u + e^- + \bar{\nu}_e \ , \qquad e^- + u \to d + \nu_e, \tag{4}$$

than normal hadronic matter does via the analogous weak nucleonic processes

$$n \to p + e^- + \bar{\nu}_e \ , \qquad \bar{e} + p \to n + \nu_e \ . \tag{5}$$

Thus neutron stars containing massless quark matter (or pion condensed matter) would, at a given age, be cooler than if they contained only normal nuclear matter. While the recent Einstein Observatory X-ray satellite measurements, which generally provide only upper bounds on neutron star surface luminosities, do indicate rather rapid cooling of the neutron stars studied, present obervations are consistent with cooling by ordinary matter.[13,14)] However, future observations, with the Exosat or Röntgen satellites for example, could, by sharpening these bounds, provide evidence of quark matter in neutron star interiors.

The early universe, as it evolved in the first few milliseconds after the big bang, followed a trajectory as shown, with the mean baryon density much less than the density of thermally excited quanta. As we know from present observations, $n_b \sim 10^{-9}\text{-}10^{-10} n_\gamma$, where n_γ is the (now 3°K) black-body photon density. The mass density, ρ, in the early universe is given by

$$\rho = \frac{\pi^2 T^4}{30\hbar^3 c^5}\, g(T) \tag{6}$$

where g(T) is the effective number of helicity states of the elementary particles present. [For $T \lesssim 1$ GeV, this number is 27 from γ, $e^\pm$, $\mu^\pm$ and gravitons, plus $\frac{7}{4}(N_\nu + 6N_f)$, where N_ν is the number of neutrino species and N_f the number of quark flavors.] Thus from solving the equations of motion for the expansion of the universe one finds that the temperature as a function of time is

$$T_{MeV} = \frac{1.56}{g(T)^{1/4}\sqrt{t_{sec}}} \approx \frac{0.50}{\sqrt{t_{sec}}} , \qquad (7)$$

where T is measured in MeV, and t in seconds. Prior to the first few microseconds, the temperature exceeded the vacuum deconfinement temperature, and hadronic matter was in the form of a quark-gluon plasma. Whether there exist present remnants of this early phase or of the transition to hadrons, such as fluctuations created in a first order phase transition,[13] remains an outstanding cosmological problem.

Central heavy ion collisions at Bevalac energies ($\lesssim$ 1 GeV/A) produce compressed excited matter, which to the extent that the internal excitation can be described by a temperature, would follow a trajectory to the upper right of the cold normal nuclear matter point on the phase diagram. It is unlikely, however, that the energy density reached in such relatively low energy collisions is sufficient to cause deconfinement of the hadrons. On the other hand, central collisions at ultra-relativistic energies, ~25-50 GeV/A, offer the possibility of exploring two regimes of the phase diagram, with production of quark-gluon matter in both. At these energies nuclei are reasonably transparent to each other so that, as shown in Fig. 3, a central collision is expected to result in two highly excited fragmentation regions containing the net baryon number of the system, one corresponding to the target and the other to the projectile nucleus; in the center-of-mass frame these regions appear as very Lorentz contracted pancakes. Similar to nucleon-nucleon collisions, these are joined together by a central region, with small net baryon number but substantial energy density. As Anishetty, Koehler

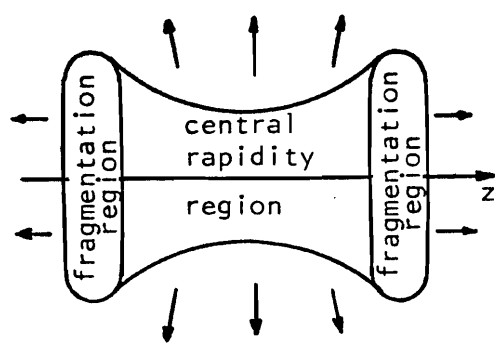

Fig. 3. Schematic picture of an ultra-relativistic central collision in the center-of-mass frame of two heavy ions, showing the two Lorentz contracted nuclear fragmentation regions, the central rapidity region, and the transverse expansion.

and McLerran[16] estimate, in the collision of two uranium nuclei at ~30-50 GeV per nucleon center-of-mass energy, the excitation energy in the nuclear fragments can easily reach ~3.5 GeV per nucleon, corresponding to an energy density $E \sim 2$ GeV/fm^3. Such a large energy density appears sufficient to produce a quark-gluon plasma, in local thermodynamic equilibrium. The region of the phase diagram reached in the fragmentation regions is indicated in Fig. 2.

The central rapidity region in such collisions can, as Bjorken[17] has recently estimated, also be excited to similar energy densities. His argument, a simple extrapolation from p-p collisions, is the following: In p-p collisions in this energy range the density of charged pions produced in the central rapidity region, per unit rapidity y, is a factor 3/2 times the observed value ~ 3 for charged pions. The energy density per unit of rapidity, assuming a final pion energy ~ 0.4 GeV is thus $d\varepsilon/dy \sim 1.8$ GeV, and increased by a factor roughly ~ A in A-A collisions.[18] Since near the central slice (z = 0, where the collision axis is along z), the velocity is $v \approx z/t$, where t is the time, the net energy density in a central A-A collision is

$$E \sim \frac{A(d\varepsilon/dy)_{pp}}{\pi R_A^2} \frac{dy}{dz} \sim \frac{0.5A^{1/3}}{t} \frac{\text{GeV}}{\text{fm}^3} , \tag{8}$$

with t measured in fm. Thus at a time of the order of a few fermi after the formation of the central region, the energy density in a U-U collision is also of the order of several GeV/fm^3, again adequate, by this time, to produce a quark-gluon plasma in thermal equilibrium. Since the baryon density is negligibly small, the matter in the central region follows a trajectory very similar to that of the early universe; in other words, ultra-relativistic central heavy ion collisions provide the possibility of studying in the laboratory the strong interaction physics that took place microseconds after the big bang.

4. Quark-Gluon Matter

From asymptotic freedom considerations, one expects quark-gluon matter to become less and less strongly interacting, logarithmically, as either the temperature or baryon density is increased, with the qcd "running coupling constant" approaching the asymptotic form

$$\alpha_s(q) = \frac{6\pi}{33-2N_f} \frac{1}{\ln(q/\Lambda)} \ . \tag{9}$$

Here q is the effective momentum scale in the matter, essentially the temperature, for very hot matter, or the Fermi momentum, for very degenerate matter; Λ, the scale parameter of qcd lies experimentally in a range around 0.2 GeV, and N_f is the number of flavors of quarks with (current) mass small compared with q.

For conditions expected in heavy ion collisions, N_f would range from 2 for u and d quarks, to 3 for u, d and s. The total number of internal quark degrees of freedom is then $12N_f$, where 12 is 2 (spin) × 3 (color) × 2 (quarks and anti-quarks); in addition the gluons each have 2 spin degrees of freedom, so that some 40 to 52 internal degrees of freedom would be excited (compared with the 4 in low energy nuclear physics). In a non-interacting quark-gluon plasma at temperature T, the energy density of the quarks is

$$E_q = \frac{7\pi^2}{20} T^4 N_f \ , \tag{10}$$

assuming equal numbers of quarks and anti-quarks, while the energy density of the gluons has the Stefan-Boltzmann value,

$$E_g = \frac{8\pi^2}{15} T^4 \ , \tag{11}$$

which is ~ 40% of the total energy. The total energy is the sum of (10), (11), and B, the difference in energy density of the perturbative and normal vacuum:

$$E = \frac{8\pi^2}{15} \left(1 + \frac{21}{32} N_f\right) T^4 + B \ . \tag{12}$$

For $N_f = 2$, the temperature is given by

$$T = 158.5(E-B)^{1/4} \text{ MeV} \tag{13}$$

where E and B are in GeV/fm^3; for $N_f = 3$ the coefficient is increased to 149 MeV. This result is useful as a first estimate of the temperatures expected in heavy ion collisions, in terms of the energy density; it implies, in particular, that because of the large number of

degrees of freedom of a quark-gluon plasma, the system remains relatively cool even at large excitation energy densities.

The number of quarks thermally excited at given energy density in hot quark matter in equilibrium is

$$n_i = \frac{9\zeta(3)}{\pi^2} T^3 = 0.14\ T_{100}^3\ (\mathrm{fm}^{-3}) \tag{14}$$

where i runs over all flavors whose mass is negligible compared with T, and in the latter form T is measured in units of 100 MeV; similarly the gluon density is

$$n_g = \frac{16\zeta(3)}{\pi^2} T^3 = 0.254\ T_{100}^3\ (\mathrm{fm}^{-3})\ . \tag{15}$$

With (12) for $N_f = 3$, the total density of quanta excited is

$$n = n_g + n_q + n_{\bar{q}} = 2.25\ (E-B)^{3/4} (\mathrm{fm}^{-3}) \tag{16}$$

where E-B is in GeV/fm^3. The large density of quanta (16) implies that they have relatively short mean free paths

$$\lambda = \frac{1}{n\sigma} = \frac{4.4}{\sigma_{mb}} (E-B)^{-3/4}\ (\mathrm{fm}), \tag{17}$$

where σ_{mb} is a mean scattering cross section in mb; and indicates that at values of E of a few GeV/fm^3, mean free paths will be small compared with the interaction volume, $\sim A^{1/3}$ fm in the transverse direction, and hence that one can use hydrodynamics to describe the initial phase of the expansion after a collision.

The total entropy density, $s = (4/3)(E-B)/T$, of matter in thermal equilibrium, with $N_f = 3$, is

$$s = 8.95\ (E-B)^{3/4} (\mathrm{fm}^{-3}) \tag{18}$$

where E-B is in GeV/fm^3. Taking E-B = 3 GeV/fm^3 we see that this entropy corresponds to one binary choice per 0.03 fm^3; the clarification of the detailed mechanism by which such entropy would be produced in the early stages of a heavy ion collision is an important question.

It is interesting to compare the above numbers with those for cold non-interacting quark matter at finite baryon density. There one has $6N_f$ Fermi seas of quarks (2 for spin, times 3 for color), with a total energy density, for $N_f = 2$, neglecting the quark masses,

$$E = 0.095\ (n_b/n_o)^{4/3}\ \left(\mathrm{GeV/fm}^3\right) + B\ . \tag{19}$$

For comparison ordinary symmetric nuclear matter has an energy density $E_o = 0.15\ \mathrm{GeV/fm}^3$. To reach energy densities of several $\mathrm{GeV/fm}^3$ in cold matter would require compressing it to $n_b \sim 10\ n_o$.

Let us turn to describe the equation of state of quark matter. The general outline of the equation of state for zero baryon number is determined by the fact that at high temperature the pressure $P = Ts-E$ of the quark-gluon plasma approaches the asymptotic value

$$P_{qg} \to \frac{1}{3}\,(E-4B) \;=\; \frac{8\pi^2}{45}\left(1 + \frac{21}{32}\,N_f\right) T^4 - B\ , \tag{20}$$

while at very low temperatures, in the confined phase, consisting essentially of pions, P is of the form (neglecting π-π interactions)

$$P_\pi = \frac{\pi^2}{15}\,T^4 b(T/m_\pi) \tag{21}$$

where b increases monotonically from 0 at $T = 0$, to unity for $T \gg m_\pi$. At a given T, the phase with the higher pressure is the favorable one, and so the deconfinement phase transition is signaled by the crossing of the two pressure curves; P/T^4 rises continuously from the value given by (21) to that given by (20). [However, since $s = \partial P/\partial T$, the equation of state P(T) has a discontinuous slope at the deconfinement transition if it is first order.]

So far in this section qcd interaction effects have been neglected; these produce deviations of the equation of state from its high temperature limiting form (20), which can be calculated by perturbative qcd. Quite generally one can show from renormalization group arguments[19)] that the pressure has the form

$$P(T) = T^4 f(\alpha_s(T)) \tag{22}$$

where $\alpha_s(T)$ is the qcd running coupling constant at temperature scale T, whose asymptotic form is given by (9). The function f can then be computed by a perturbation expansion in α_s (reviewed in Ref. 2). Unfortunately this approach is of limited validity since non-perturbative effects appear to play a large role. A hint of trouble can be seen from the first order result for the entropy density (N_f=3)

$$s = \frac{19\pi^2}{9} (1- \frac{54}{19\pi} \alpha_s(T) + \ldots) T^4 \tag{23}$$

which turns negative for $\alpha_s > 1.1$.

The presently most promising approach to calculating the free energy of quark-gluon plasmas (and the properties of qcd more generally) is by means of Monte-Carlo studies of qcd on a lattice. Calculations by McLerran and Svetitsky,[20] Kuti et al.,[21] and Engels et al.,[22] for SU(2) lattice gauge theory at finite temperature, without quarks, real or virtual, indicate a deconfinement transition at a critical temperature $T_{dec} \sim$ 160-200 MeV; similar calculations in SU(3) by Kajantie et al.,[23] and Montvay and Pietarinen[24] yield T_{dec} also on the order of 200 MeV. In lattice gauge calculations without fermions the transition is characterized by the interaction energy between two test quarks changing from linear at large separation in the confining phase below T_{dec}, to Coulomb-like above T_{dec} in the deconfined phases. For SU(2) the transition is expected to be similar to the second order transition of the simple Ising model. However once quarks are included, one no longer can have a linear force law due to the possibility of creating $q\bar{q}$ pairs in the vacuum, "breaking the string," and the sharp transition should be either washed out as in an Ising model in non-zero magnetic field, or become first order.

Notable has been the work of the Bielefeld group[8,22] in calculating the finite temperature equation of state of a quark-gluon plasma, at zero baryon density. Figure 4 shows their results for the ratio of the energy density of the SU(3) system, with N_f= 2, to the Stefan-Boltzmann value [E-B in Eq. (12)]; we see clearly how the calculated values approach this T^4 limit. The temperatures in this figure are measured in units of Λ_L, the lattice scale parameter, which is ~ 2 MeV for SU(3), and ~ 4 MeV for SU(2). As they also show, at the low temperature end, below T_c (not shown in Fig. 4), the thermo-

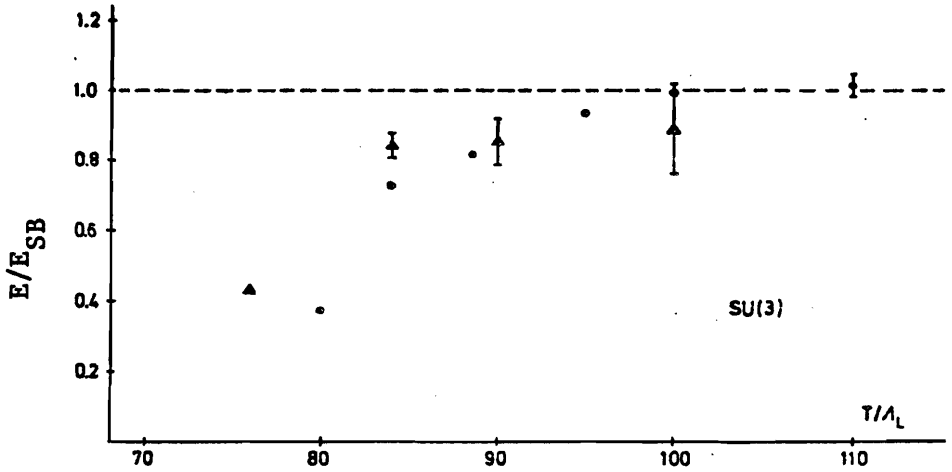

Fig. 4. Ratio of the energy density of an SU(3) quark-gluon plasma ($N_f = 2$) to the T^4 Stefan-Boltzmann limit, as a function of temperature, from the Monte Carlo calculations of Engels et al.[8)]

dynamics of the pure Yang-Mills SU(2) system (no fermions) is similar to that of a gas of "glueballs" of mass $m_g \sim 3\text{-}4\ T_c$, with an exponentially increasing density of resonance states. These calculations demonstrate the potential utility of lattice gauge theory for deriving a complete equation of state over the entire span from the low temperature hadronic limit to the high temperature asymptotically free quark-gluon limit.

In addition to Monte Carlo calculations of the deconfinement temperature, the deconfinement curve (Fig. 2) has been studied by several other methods, which we briefly survey. [For a more detailed review see Ref. 2.] As a first estimate of the limits of the region where it makes sense to describe matter in terms of distinct hadrons, we note that the density of cold nuclear matter at which the nucleons begin to overlap substantially is $\sim (4\pi r_n^3/3)^{-1} \equiv n_h$, where r_n is the nucleon radius; for r_n between 1 and 0.4 fm., n_h is between 1.4 n_o and 23 n_o. Similarly the temperature (for $n_b = 0$) at which thermally created pions begin to fill space is determined by $n_\pi(T) \sim (4\pi r_\pi^3/3)^{-1}$, where $n_\pi(T) \simeq 3\zeta(3)T^3/\pi^2$ (for $T \gg m_\pi$) is the density of thermally created pions; for a nominal pion radius $r_\pi \sim 0.6$ fm, this temperature is $\sim 2m_\pi$. [A sophisticated version of these estimates using a self-consistent statistical bootstrap spectrum of hadronic excitations is given in Ref. 25.] Beyond these ranges the basic degrees of freedom of the system are quark-like rather than hadronic.

Starting from the quark-gluon phase, one may expect, with Kajantie and Miettinen,[26] a transition when the density of quarks plus anti-quarks present falls to a value on the order of that in hadrons, $\sim 3n_h$; neglecting interactions and the baryon chemical potential in evaluating the density of anti-quarks present, one finds that this criterion yields a deconfinement curve

$$T_{dec}(n_b) = 276(n_h - n_b)^{1/3} N_f^{-1/3} \ (\mathrm{MeV}) \tag{24}$$

where N_f is the number of quark flavors whose anti-particles are excited at temperature $T_{dec}(n_b)$, and the densities are measured in fm^{-3}. This result at $n_b = 0$ is in qualitative agreement with the other estimates above. Its disadvantage is that it does not take into account effects of thermal gluons, which can contain up to nearly half the energy of the system. A possibly better phenomenological criterion for the transition is to say that in quark matter at finite T, confinement takes place when the energy density falls to a certain critical value E_c. If we neglect interaction effects, Eq. (13) then relates $T_{dec}(n_b = 0)$ to E_c, while Eq. (19) implies $n_{dec} \simeq (0.94\ fm^{-3})E_o^{3/4}$, where E_o is E_c-B in units of GeV/fm^3. The value of E_c can be estimated from the point where the ground state energies densities, as a function of n_b, of nuclear and quark matter cross; one finds[2] $E_c \sim 1.3$-$1.8\ GeV/fm^3$. [Such comparison of ground state energies, or free energies more generally at finite T, itself provides another way of estimating the deconfinement transition.] The corresponding T_{dec} then agrees well with previous estimates. The general conclusion is that for small n_b, deconfinement should set in at temperatures $\lesssim 200$ MeV, with corresponding energy densities $\lesssim 2\ GeV/fm^3$, values seemingly attainable in ultra-relativistic heavy ion collisions.

5. The Chiral Transition

One of the more intriguing results to emerge from recent lattice gauge theory calculations has been the demonstration by the Illinois group that as the temperature of the vacuum is raised, the deconfinement transition, at T_{dec}, somewhat precedes the transition, at T_{ch}, in which chiral symmetry is restored,[7] a result which also appeared in

the calculations of the Bielefeld group.[8] Such a possibility was suggested earlier by Pisarski,[27] who gave arguments that T_{ch} is at least T_{dec}, but not too much greater; and independently by Shuryak,[28] starting from the observation that the qcd interaction scale parameter Λ, and constituent quark masses provide two separate hadronic length scales, the first related to confinement and the second to chiral symmetry breaking.

The Illinois calculation is done by imposing a small quark mass term $m\bar{q}q$ in the Lagrangian to break chiral symmetry explicitly, and studying the limit of $\langle\bar{q}q\rangle$ as m, the imposed quark mass, goes to zero. This method is analogous to that used to study spontaneous magnetization by applying a weak external field to define a direction for the magnetization, and taking the limit of vanishing external field. Internal quark loops are neglected. As we see in Fig. 5 from Ref. 7, $\langle\bar{q}q\rangle$ becomes non-vanishing in this limit provided the coupling g^2 is sufficiently large, or equivalently, that the temperature is sufficiently small, thus indicating spontaneous breaking of chiral symmetry below a critical T_{ch}, with the quarks becoming massive. For SU(2) lattice gauge theory, the calculations show that T_{ch} is on the order of 10-20% higher than T_{dec}, and definitely distinct. To test the independence of the deconfinement and chiral transitions, the chiral properties of quarks in the adjoint representation of the gauge group were also studied; on the one hand, quarks in this representation can be perfectly screened by gluons, so they are not confined in pairs [in SU(2)], while on the other

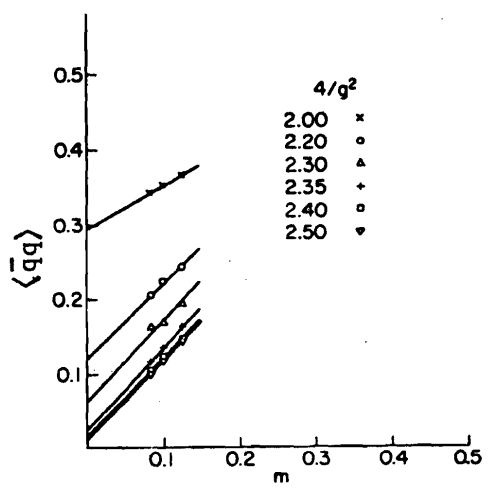

Fig. 5. Chiral order parameter $\langle\bar{q}q\rangle$ as a function of the bare quark mass m, for various coupling strengths. The highest curve corresponds to the lowest temperature. [From Ref. 7.]

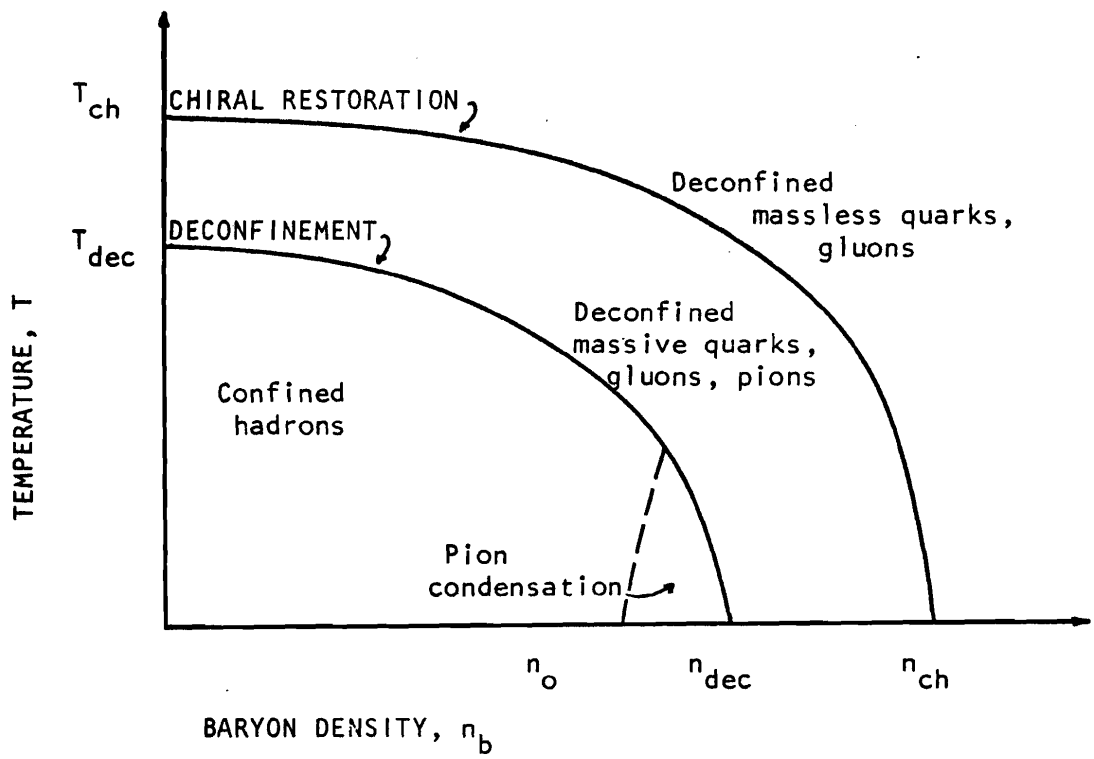

Fig. 6. The three phases of strongly interacting matter: confined, deconfined with spontaneously broken chiral symmetry, and deconfined with restored chiral symmetry.

hand T_{ch} is substantially raised. Generally, the chiral transition appears to depend more strongly on the gauge group than the deconfinement transition.

These calculations indicate that the vacuum, and indeed matter at finite baryon density, should exhibit <u>three</u> distinct phases: at low T and n_b, the vacuum is confining, and matter is in the form of hadrons, including, as the Goldstone bosons of the broken chiral symmetry, (nearly) massless pions. As before, at very high T or n_b, matter is composed of deconfined massless (for u and d) quarks, and chiral symmetry is restored. However, there should now be an intermediate region just beyond the deconfinement curve, yet another state of matter, consisting of deconfined quark-gluon plasma; with chiral symmetry is still spontaneously broken. The quarks in this region are therefore all massive; perhaps surprisingly, the system still has massless pionic excitations. [One should note that were the situation reversed, with chiral symmetry restoration preceding deconfinement,

the intermediate phase would correspond to Lee-Wick abnormal matter,[29] in which one has massless nucleons.] The three phases are illustrated in the new phase diagram, Fig. 6, a more accurate version, one hopes, of the old phase diagram, Fig. 2. [As we have discussed earlier, the deconfinement transition may not be perfectly sharp; since chiral symmetry is not exact, the chiral transition should also be slightly washed out.] The actual locations of the two transition curves cannot be given as yet with any precision; in particular the possibility is open that one may only reach the intermediate deconfined region in heavy ion collisions, and not the fully chirally symmetric phase.

More generally, there is a present need for reliable predictions of the properties of the intermediate phase, such as the equation of state, the variation of quark masses with T and n_b, and the order of the chiral transition. Where the transition is second order, the u and d quark masses should tend smoothly to zero at the transition. As a first approach, we have recently studied a model description of this phase,[30] based on the Friedberg-Lee soliton model of hadrons,[31,32] and find, for example, that the transition at high baryon density and low temperature can be either first or second order, dependent on the choice of parameters, but always second order at low n_b and high T.

A natural question to ask about the massive quark phase, since it has pionic excitations, is whether it might be pion condensed in certain regions of T and n_b, i.e., have a macroscopic condensed pion field in its state of lowest free energy. The pion condensation problem is particularly simple in this phase; the pseudovector πqq coupling constant, $g_{\pi qq}/2m_q$ is equal, by the Goldberger-Treiman relation, to g_A/f_π, where m_q is the quark mass, g_A is the quark axial current renormalization, and f_π is the pion decay constant. Since for point quarks, $g_A = 1$, the πqq coupling constant is just $1/f_\pi$. Also the system does not have the complications of coupling to nucleon resonances, Δ, as in the nuclear matter pion condensation problem. However, as one finds in the analogous nucleon problem without Δ's,[19] condensation, at least at $T = 0$, depends essentially on g_A being > 1, which suggests that the massive quark phase should not undergo pion condensation.

A second question is whether this phase in neutron star matter should exhibit the enhanced cooling, via the reactions (4), expected for the massless quark phase. For degenerate matter in beta equilibrium under the reactions (4), with $m_u \simeq m_d$, one finds readily that the down quark Fermi momentum exceeds that of the electron plus up quark. Consequently the decay $d \to e^- + u + \bar{\nu}_e$ cannot conserve momentum if it conserves energy ($E_\nu \sim T$), and hence the cooling would proceed with the same T dependence and smaller magnitude as the nucleon reactions (5). In other words, the massive quark phase would not be readily detectable in neutron stars through its cooling properties.

6. Conclusion

While considerable theoretical work still needs to be done to understand the properties and parameters of the deconfined phases of matter, and the transition from hadronic to quark-gluon matter, the broad outlines of this new phase of matter are already clear. The prospects for producing quark-gluon plasmas in ultra-relativistic heavy ion collisions are very promising. Not only would production and detection of such a basic new state, to be discussed in detail later in these proceedings, be a remarkable achievement in itself, but by enabling one to study qcd in extended deconfined matter over large distance scales it should reveal fundamental aspects of qcd and confinement unattainable from single hadron physics.

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