Fub-85/27-T

The Physics of the Quark-Gluon Plasma and How It Might Be Produced in Hadronic Collisions

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The simplest systems studied in physics are those consisting of one, two, or infinite numbers of particles. The symmetries and dynamics of particle interactions may simply manifest themselves in systems such as an isolated electron, a hydrogen atom or a magnet. The dynamics of the motion of an isolated electron probes the Lorentz and translation invariance of the vacuum, and is manifest in conservation laws which describe the electron's motion. The non-existence of isolated quarks is the most compelling evidence for the confining nature of the theory which describes strong interactions, Quantum Chromodynamics (QCD). In two particle systems such as the hydrogen atom, the dynamics of particle interactions are isolated and may be simply studied. In multi-particle systems such as magnets, systems which are described by Ising models and their ilk, the symmetries of particle interactions may arise in the phase structure of the systems as symmetries are broken or realized. In a magnet, the fundamental spin-spin interactions which magnetize a solid are invariant under a spatial rotation of all the spins in the magnet. At high temperatures, this symmetry is demonstrated by the inability of a magnetic solid to sustain a magnetic field. At low temperatures, the magnet is magnetized, and the symmetry of rotational invariance of the fundamental spin-spin interactions is broken, since a preferred axis of magnetization is chosen by the magnet. The magnetizationdemagnetization phase transition simultaneously demonstrates the rotational invariance of spin-spin interactions, and the associated intrinsic magnetization of atoms with net spin. For infinite particle systems the dynamics simplifies since it may be studied by powerful statistical mechanical methods. The dynamics of few particle systems containing three or more particles but not infinite numbers of particles, are intrinsically complicated and extracting the fundamental particle symmetries and dynamics is difficult.

Much progress in understanding strong interaction physics has arisen from studying systems which are approximately composed of two particles. For example, in a high energy electron-positron collision, a high energy quark-antiquark pair is formed. The quark-antiquark pair have a very large net energy. In the reference frame in which they are formed at rest, the large total energy must be given to the quark-antiquark system which has zero net momentum. This is only possible if the quark-antiquark pair travel away from the region where they were created with equal and opposite momentum of large absolute value. If either the quark or anti-quark of the pair form only a few more particles, many of which have momentum close to that of either the quark or antiquark, then the production of the quark-antiquark pair produces a spray of particles with a distinctive topology. The particles move largely along an axis parallel to the momentum of either of the quark or anti-quark. Two back to back jets of particles are produced, and the existence of the jets provides a striking confirmation of the picture that only a single high energy quark-antiquark pair was produced when the high energy matter formed.

If the quark-antiquark pair was to emit many particles of momentum not parallel to that of the quark or anti-quark after its formation, and these many particles had more or less random momentum, then the large momenta of the formed quark anti-quark pair would be degraded. The direction of motion of the quark would no longer be opposite to that of the anti-quark, and the produced particles no longer necessarily appear as a jet. As we shall discuss in the next paragraph, this does not occur if the quark-antiquark pair have a very high energy, that is, an energy large compared to the energy stored in the rest mass of the proton.

The simplicity of jet formation in high energy electron-positron annihilation arises from the asymptotic freedom of strong interactions. At high energies, or short distances, the strong interactions become weak. Strong interactions in general change particle number, and the creation of particles is therefore suppressed at high energies and short distances. This fact allows for a systematic expansion in powers of the strong interaction coupling constant when computing the properties of jets. Such an expansion is to a large extent just an expansion in the number of particles which might degrade the momentum of the quark-antiquark pair and spoil the jettiness of the quark-antiquark jet. Of course, as the quark-antiquark pair initially produced in the collision separate to larger distances, the strong interactions become strong and multi-particle degrees of freedom become important. A remarkable feature of QCD is that when these interactions begin to become strong, for many of the properties of jet production the basic jettiness of the jet is not destroyed. This is because the particles which are produced by the quark and anti-quark in the jet have only a small momentum perpendicular to the direction of motion of the quark and anti-quark.

Multiparticle dynamics and strong interactions, when interactions are truly strong, seem to be locked together. This is an almost inevitable consequence of the particle number non-conserving nature of the strong interactions. Even the ground state vacuum of quantum chromodynamics is a state consisting of an infinite number virtual quarks, antiquarks and gluons. The quarks, antiquarks, and gluons are the elementary excitations of the theory at high energy and short distances. When interactions are strong, these excitations are no longer elementary, and must travel through and polarize a vacuum which consists of an infinite number of strongly interacting elementary excitations. Statistical mechanics techniques seem ideally suited for analyzing such systems.

As suggested by these considerations, QCD studies have a dual nature. One goal of such studies is to verify that QCD in fact describes strong interaction processes. For this purpose, studies of few particle systems should be adequate. When few particle systems are isolated, the elementary interactions between the particles may be probed. For example, when a three jet system is produced in electron-positron annihilation, the basic interaction where a quark or anti-quark produces a quark or anti-quark plus a gluon may be measured. This production process converts the two jet system of a quark and anti-quark into a three jet system of a quark, anti-quark and gluon. At high energies, this process is infrequent since the particle number changing strong interactions are weak. Such processes are nevertheless relatively easy to observe, and because the particle number changing strong interactions are weak, may be isolated from four jet processes.

Another goal is of QCD studies is to understand the dynamics of QCD when interactions are truly strong. These studies may lead to fundamental understanding of the nature of confinement, that is, the non-existence of isolated colored particles in nature, and the nature of chiral symmetry breaking, that is why the proton is so massive compared to the pion. Chiral symmetry breaking is particularly mysterious, since chiral symmetry is believed to be a symmetry of QCD which is analogous to rotational invariance for magnetic systems. If the symmetry was realized in nature, as the rotational symmetry is realized for a demagnitized high temperature magnet, then protons would either have zero mass, or would be accompanied by mirror particles of equal mass. In nature, strongly interacting mesons have masses small compared to the proton, and it is difficult to imagine the proton as a light particle. There are also no equal mass mirror particles for the proton. The chiral symmetry is believed to be broken.

Since the study of strong interaction dynamics demands the study of multiparticle systems, we should ask what are the simplest multiparticle systems to study. The vacuum or ground state of QCD is perhaps the simplest such system, but suffers from the fact that there are no adjustable parameters which charachterize the vacuum. The vacuum has nothing in it, so what can be adjusted? Introducing strongly interacting particles strongly polarizes the vacuum, and the system becomes even more complicated. With no adjustable parameters, it is difficult to imagine experimentally probing the vacuum. Also, with adjustable parameters, systems often become simple for some ranges of the parameters. For these ranges, theoretical studies simplify, and some insight may be gained about the system as the parameters vary into a region where the dynamics becomes more complicated.

The simplest set of systems characterized by a minimal set of adjustable parameters are finite temperature and baryon number density hadronic systems. Two parameters characterize these systems, the temperature T and the baryon number density, ρ_B. Baryon number density is the number of baryons minus antibaryons, that is, protons plus neutrons (plus other baryons) minus antiprotons and antineutrons (and other antibaryons), per unit volume. Baryon number is conserved by the strong interactions. In the limit that $T \to 0$ and $\rho_B \to 0$, these systems become the vacuum. Finite temperature and density systems provide the simplest generalizations of the vacuum. A remarkable property of these systems is that the mathematics of the study of elementary excitations is in exact analogy with that of the study of the elementary excitations of the vacuum. The elementary excitations of the vacuum are elementary particles such as gluons and quarks. For high energy density systems, the description of these elementary excitations simplify. At high energy density, the average particle separation is small, so that the effective strength of interactions is weak. For many purposes, the elementary excitations may be approximated as free quarks and gluons in a finite temperature and baryon number density ensemble.

We shall mean by high energy density matter, matter with energy densities greater than or of the order of the energy density of matter inside of nuclei. By ordinary standards, this energy density is tremendous, but for our purposes it is the minimal value at which the physics of QCD becomes interesting. The energy density of nuclear matter is $\rho \sim m_p c^2 A / (4/3 \pi R^3)$ where m_p is the proton mass, A is the nuclear baryon number, and R is the nuclear radius. Since $R \sim 1.1\ A^{1/3}$ Fm. where 1 Fm. = 10^{-13} cm, this energy density is $\rho \sim .15$ Gev/Fm^3. At energy densities typical of that inside of nuclei, quarks and gluons are very probably confined, on the average, inside of hadrons such as protons and neutrons, or pions. At much higher energy density, when all of these hadrons are squeezed so tightly together that on the average they all overlap, the system may become an unconfined plasma of quarks and gluon which are free to roam the system. The energy scale at which the hadrons begin to overlap is roughly the energy density of matter inside of a proton, $\rho \sim m_p c^2/(4/3\ \pi R^3)$. If we use the r.m.s. radius of the proton as measured in electron scattering, $R \sim .8$ Fm., this energy density is $\rho \sim .5$ Gev/Fm^3, an energy density not so large compared to that of nuclear matter. If nuclei are squeezed to perhaps only an order of magnitude higher energy density, perhaps a transition to a quark-gluon plasma results. At energy densities much higher than that inside a proton, the matter is at such a high density that interactions are unimportant, and the system, for many purposes, becomes an almost ideal gas of quarks and gluons.

High energy density hadronic matter occurs in a variety of physically interesting situations. Cold, baryon rich hadronic matter of almost infinite extent occurs in the cores of neutron stars at energy and baryon number densities of maybe 10-20 times that of ordinary nuclear matter. This matter is quite cold compared to the kinetic energy scale of a typical baryon in the star. Typical neutron star temperatures correspond to an energy scale of Kev's where the typical kinetic energy of a baryon at nuclear matter energy density is ten to twenty Mev. The

Fermi energy which charachterizes the average baryon kinetic energy in this high density matter arises from the fact that if the system has its baryon number increased, the Fermi exclusion principle requires that the newly added baryons must occupy presently unoccupied states of higher energy than those which are presently occupied. Increasing the baryon number density, therefore increases the typical particle kinetic energy and therefore results in a Fermi energy. Since the natural time scale for strongly interacting systems is the time it takes light to travel a Fermi, 10^{-23} sec., and since neutron stars have macroscopic lifetimes, the cores of neutron stars almost certainly contain matter which is in local thermal equilibrium and may be well described using equilibrium thermodynamics.

Another example of a high energy density system is provided by the big bang. In the big bang, energy densities may have existed up to at least those where quantum gravity becomes important, $\rho \sim 10^{70-80}$ times that of nuclear matter. When the energy density was that of ordinary nuclear matter, the universe was expanding with a charachteristic time scale of $t \sim 10^{-6}$ sec. At this time the matter was hot, with only a small baryon number density. Put another way, the typical excitation energy of particles was so large, and the number of states available for particles to occupy was so large, that only occasionally would two particles try to occupy the same state. At temperatures where the system may be described as a quark-gluon plasma, in addition to the baryon number carrying quarks, a large amount of energy was carried by gluons. This situation is unlike the neutron star case where almost all of the energy resides in quarks. Below the temperature at which hadrons are confined, a large amount of energy is carried by mesons, such as the pions which carry no baryon number, as well as baryons. For the neutron stars at densities below that of confinement, the energy is almost entirely carried by baryons. The charachteristic time scale of expansion is orders of magnitude larger than the natural time scale for strongly interactions, and the universe was probably to a very good approximation in local thermal equilibrium.

Another place where hot, baryon rich matter might be produced is in ultra-relativistic heavy ion collisions. The result of such a collision is shown in Fig. 1.[1-2] Here the two Lorentz contracted nuclei pass through one another leaving hot hadronic matter in their wake. In the region of the Lorentz contracted nuclei, hot, baryon rich matter presumably exists. This picture of ultra-relativistic nuclear collisions will be explained in more detail later. This situation is much different than that at very low energies where the nuclei are almost impenetrable, and is a consequence of the Lorentz time dilation of hadronic interactions of ultra-relativistic particles.

Reasonable estimates of the energy densities achieved in these collisions are 10-400 times that of ordinary nuclear matter. Particles with velocities close to that of the ultra-relativistic colliding nuclei come from a fluid with a velocity close to the velocity of the nuclei, and the fluid as viewed in a reference frame moving with the fluid, has a high energy density. This 'nuclear fragmentation region' fluid has high baryon number density, and the matter has properties similar to that produced in neutron stars. The matter produced with velocities low compared to that of the colliding nuclei, also have a high energy density as measured in a comoving reference frame, and this energy density is not so baryon rich. The matter produced in this 'central region', that is with velocities central between those of the two colliding nuclei, is therefore analogous to that in the early universe. The time scales for expansion of such matter are perhaps only a few times to 100 times the natural time scale for hadronic interactions, and although the system may be to a fair approximation in local thermal equilibrium, non-equilibrium processes will also be important for the dynamics.

The description of hadronic matter by QCD, ignoring interactions other than strong interactions, is adequate for a range of energy densities less than those for which heavy W and

Z bosons, the particles responsible for weak interactions, form an important component of the matter. Electromagnetic effects must also be included at all energy densities and are more important than strong interactions for describing dynamics at energy densities much less than those of nuclear matter. At low energy densities, particles are well separated. Since strong interactions are of short range, the long range electromagnetic interactions, which are weak at distance scales smaller than a Fermi, begin to dominate. At extremely high energy densities, $\rho \sim 10^{60-80}$ times that of nuclear matter, particles with the degrees of freedom associated with grand unification, and eventually even quantum gravity, may play a dominant role. At any energy density at which new degrees of freedom play an important role, a variety of phase changes may take place. Such energy densities are shown in Table 1. We shall be interested in matter in the range of energy densities 10^{10}Gev/Fm^3 $>> \rho \gtrsim 1$ Gev/Fm^3, where QCD adequately and simply describes matter. At these densities, electromagnetic interactions contribute to the dynamics, but play a sub-dominant role since their interactions are weak compared to the strong interactions.

The dynamics of hadronic matter as described by QCD is quite different for some ranges of energy densities. At energy densities which are very large or very small compared to some natural scale, we might expect the dynamics of the matter to simplify. At energy densities close to some natural scale we might expect that the dynamics would be complicated, but perhaps more interesting than in the asymptotic regimes. At such energy densities, a variety of phase transitions, or at least dramatic qualitative changes in the properties of hadronic matter, might occur, and their existence might reflect symmetries of the QCD interactions. At energy densities which are large compared to this natural scale, $\rho \sim .5$ Gev/Fm^3, hadronic matter is a plasma of quarks and gluons unconfined into hadrons. At such very high energy densities, the average separation between the constituents of the quark-gluon plasma may become $d << 1$ Fm. The strength of the QCD interactions between these closely packed quarks and gluons is weak due to asymptotic freedom, since QCD interactions become weak at short distances. Of course, there may be some long range interactions which are not shielded in the plasma and for such long range interactions the strength is not small. Such long range interactions do not however significantly affect the bulk properties of the plasma, but might significantly alter long distance correlations between particles in the plasma. Since the interactions which affect the average bulk properties of the plasma are weak, they provide small corrections to the kinetic energies of the quarks and gluons when bulk quantities such as the energy density and the pressure are computed. Hadronic matter therefore becomes an almost ideal gas at very high energy densities.

At low energy densities, hadronic matter is composed of widely separated hadrons. Since the range of strong interactions is finite, these widely separated hadrons will, on the average, interact weakly. Hadronic matter becomes an ideal hadron gas at these low energy densities, except for corrections due to electromagnetic interactions. At extremely low temperature, these electromagnetic interactions may eventually dominate over kinetic energies of particles, and solids or liquids may form. For some range of temperatures, we nevertheless expect that the hadronic plasma will be adequately described as an ideal gas.

The properties of hadronic matter simplify at very high and very low energy densities, when hadronic matter becomes an ideal gas. The degrees of freedom of these ideal gasses are of course quite different. For a zero baryon number density system, the low temperature system is a pion gas with three degrees of freedom. At very high temperatures there are colored spin 1/2 quarks and antiquarks with three colors corresponding to 12 degrees of freedom plus eight colored helicity 1 gluons, with two allowed polarizations, corresponding to 16 degrees of freedom. For intermediate energy densities, hadronic matter is interpolating between these degrees of freedom. The number of degrees of freedom change by approximately an order of magnitude in going from a pion gas to a quark-gluon plasma. This huge change is reflected in the magnitude of the change in the energy density of the matter. Including the rest masses of pions, the energy density of the quark-gluon plasma, just above the confinement-deconfinement phase transition changes by an order of magnitude. Before the development of QCD, it was believed that at some temperature, the energy density of hadronic matter diverged, and there was therefore a limiting temperature for hadronic matter. Within QCD, this divergence is replaced by a large change, and temperatures in excess of the Hagedorn limiting temperatures are realized.

At the energy densities which interpolate between a pion gas and a quark-gluon plasma, a variety of phase transitions may occur. As the quarks and gluons become liberated from hadrons and the degrees of freedom of hadronic matter dramatically change, a confinement-deconfinement phase transition may appear. The singular confining force between quarks and gluons may also rapidly change. In addition to the confinement-deconfinement phase transition, a variety of other transitions may independently or simultaneously arise. An example is the chiral symmetry restoration phase transition. Chiral symmetry is a symmetry of the QCD interactions when the up and down quarks are approximated as massless. Since the up and down quark masses are believed to be very small compared to the typical mass scales of QCD, chiral symmetry is a good approximate of actual QCD interactions. A consequence of chiral symmetry is that either all baryons, such as the proton and neutron, are massless, or there are mirror particles to the proton and neutron which are degenerate in mass. Neither of these two possibilities seems to be realized for the baryons observed in nature, and the chiral symmetry is believed to be dynamically broken. This situation is analogous to that of magnetic systems. Even though the basic spin-spin interactions are invariant under a spatial rotation which rotates all the spins in a magnet, the direction of magnetization of a magnet chooses a preferred spatial orientation, and the properties of a magnet are not rotationally invariant. It is believed that chiral symmetry is broken in the same way. As a consequence of this symmetry breaking, the pion is believed to acquire a zero mass, which is a fair approximation with what is observed. In analogy to the magnetic spin system, perhaps chiral symmetry, like rotational invariance may be restored at high temperatures, and baryons might become effectively massless and the pion effectively massive.

The transition between the degrees of freedom of ordinary hadronic matter and that of a quark-gluon plasma may be shown to be a symmetry breaking phase transition, like that of a spin system, when QCD is approximated by taking all quark masses to be very large compared to the mass scale set by confinement.[2-3] (The mass scale of confinement is determined by asking what is the mass of a particle whose Compton wavelength, $\lambda = h/mc$, corresponds to the length at which hadronic interactions are strong.) In this approximation, confinement arises as a consequence of symmetry breaking, and deconfinement as a consequence of symmetry realization.

To probe the gluon plasma (in the absence of dynamical quarks), heavy test quarks may be inserted into the plasma. The free energy, $F(\vec{r})$, of a test quark at position $\vec{r}$ in a system with temperature $\beta = T^{-1}$ may be found by computing the avearge value of a spatially local operator $L(\vec{r})$. It is possible to show that the free energy of an isolated test quark at the position $\vec{r}$ is

$$\langle L(\vec{r}) \rangle = e^{-\beta F(\vec{r})} = \begin{cases} 0 & \text{Confinement} \\ \text{Finite} & \text{Deconfinement} \end{cases} \tag{1}$$

The operator L is analogous to the magnetization of a spin system. Its average value is zero in the confined phase, where the free energy of an isolated test quark is infinite, in analogy to the demagnetized phase of a spin system. In the deconfined phase, the free energy is finite, corresponding to a magnetized phase.

In this analogy to a spin system, it turns out that the temperature of the analogous spin system is not the true temperature, but is given by the strength of gluon interactions,

$$T_{Eff} = g^2 \tag{2}$$

where g is the gluon coupling strength. The effective temperature depends on the physical temperature since g^2 depends upon T. Recall that at high temperatures, the coupling strength decreases due to asymptotic freedom. Eq. 2 predicts that as the the physical temperature increases, the effective temperature decreases and vice-versa.

The operator L is analogous to spin variable. For a theory with N^2-1 gluons (N = 3 for QCD) the QCD equations of motion are invariant under a global transformation which flips L:

$$L(\vec{r}) \rightarrow e^{\frac{2\pi i j}{N}} \; L(\vec{r}) \tag{3}$$

where j is an integer. This Z_3 spin flip symmetry, which we shall call the center symmetry, guarantees that <L> = 0 if the symmetry is realized. This symmetry is analogous to rotational symmetry, where all of the spins would be rotated through a continuous real angle. In QCD, the angles are discrete, and imaginary.

At very high effective temperatures of the Z_3 spin system, the spin system should be disordered, <L> = 0, and the system confines. As the effective temperature decreases, there may be a critical temperature for which the spin system orders, <L> ≠ 0, and the system deconfines. It can be shown that if the effective spin system exhibits a first order phase transition, then this will also be true for the QCD finite temperature theory. For Z_3 spin systems, corresponding to QCD, the phase transition is predicted to be first order. Since effective temperature and physical temperature are inversely correlated, low physical temperatures correspond to symmetry restoration and confinement, while high physical temperatures correspond to symmetry breaking and deconfinement.

This picture of the confinement-deconfinement phase transition for a gluon plasma may be studied using lattice Monte-Carlo methods.[2,4] The latticization starts by discretizing space-time in a finite four volume, and physical quantities are extracted in the zero lattice spacing infinite three volume limit. The fourth length of the four volume may be shown to be $\beta = 1/T$. Using a discretization of coordinates is a familiar technique for solving partial differential equations and is of great use for numerically solving QCD. The phase structure at finite temperature arises in the zero lattice spacing, infinite three volume limit as the length of the fourth dimension of the box, $1/T$, is varied. This length might be varied by changing the coupling, g^2, since temperature and coupling are related through asymptotic freedom.

Once QCD is discretized and placed on a four dimensional lattice, it may be studied by lattice Monte-Carlo simulation techniques. Various gluon field configurations which give contributions to the partition function may be selectively sampled by an algorithm which finds those contributions which dominate the sum over configurations. The error in this selective sampling is approximately the inverse square root of the independently sampled gluon field configurations.

Results of a Monte-Carlo simulation of $\langle L \rangle$ are generically shown in Fig. 2. For $N = 3$, corresponding to QCD, there is a sharp break at the critical temperature corresponding to a first order phase transition, and actual Monte-Carlo computations yield curves for $\langle L \rangle$ similar to that shown. The specific heat and energy density as a function of temperature have also been computed. At high temperatures, these quantities approach ideal gas values. The energy density jumps rapidly at T_c, and a first order phase transition is strongly indicated. The latent heat of this phase transition appears to be 1-2 Gev/Fm^3, with uncertainties of undetermined magnitude arising from the small size of lattices used in the computations, but which might be as large as an order of magnitude. The critical temperature is determined to be $T_c \sim 200$ Mev with uncertainty of perhaps a factor of 2.

When dynamical quarks are included in finite temperature QCD computations, the confinement-deconfinement transition might dissappear. To understand how this might happen, consider the free energy of an isolated test quark. In a system with dynamical quarks, this free energy is always finite since the test quark may always form finite energy bound states with the dynamical quarks. The operator L is no longer an order parameter, since it always has a finite expectation value. In the analogy with a spin system, the dynamical quarks are analogous to an external magnetic field since their presence guaranties that the magnetization $\langle L \rangle$ is always finite. Since the effects of dynamical quarks dissappear as the quark mass becomes infinite, the external magnetic field goes to zero in this limit.

In spin systems such as the Ising model, the presence of an external magnetic of arbitrarily small strength can destroy the magnetization phase transition. The case for QCD is somewhat more complicated since in the absence of dynamical quarks, the confinement-deconfinement phase transition is first order, where for the Ising model the phase transition in the absence of an external magnetic field is second order. A first order phase transition involves a discontinuous change in the properties of a system, and a small perturbation may only continuously deform this change. Only if the effects of quarks become sufficiently large may the phase transition be removed. We therefore conclude that if the quark masses are sufficiently small and if there are a sufficiently large number of quark flavors, then the phase transition might be removed.

These arguments are complicated by the possibility of a chiral symmetry restoration phase transition. Since chiral invariance is a symmetry of QCD only when the up and down quark masses are zero, this symmetry is only approximate in nature since these masses are small but finite. A simple model for chiral symmetry breaking is the σ model where a spin zero, even parity and even charge conjugation σ meson condenses in the vacuum and induces chiral symmetry breaking. In QCD, this σ meson is presumably a bound quark-antiquark pair. As quarks and gluons begin to become deconfined from hadrons, the σ mesons become ionized. Since fermions do not condense by themselves, and since there are no longer any σ mesons to condense, chiral symmetry breaking may end.

Since chiral symmetry breaking generates dynamical masses for baryons such as quarks, and since the strength of the effective magnetic field which might destroy the confinement-deconfinement phase transition increases with decreasing quark mass, the confinement-deconfinement phase transition and the chiral phase transition are intimately related. In the approximation where the strange quark mass, as well as that of the up and down quark, is taken to be zero, it is possible to show that the chiral symmetry restoration phase transition is first order. As the up, down, and strange quarks acquire small but finite masses, this first order phase transition should not dissappear.

A possible phase diagram for the chiral and confinement phase transitions is shown in Fig. 3. The upper and bottom horizontal axis are the temperature while the vertical axis is $e^{-1/m}$. This latter variable approaches 1 for infinite mass quarks, and approaches 0 for zero mass quarks. The upper T axis therefore corresponds to a theory in the absence of dynamical quarks. Along this axis there is a first order phase transition. As the mass decreases from infinity, a line of first order phase transitions may evolve in the T, $e^{-1/m}$ plane. This line might terminate in a first order phase transition as shown in Fig. 3a or it might evolve all the way to zero T as in Fig. 3b. Along the bottom T axis, the quark mass is zero. There should be a chiral phase transition along this axis which may be first or higher order. If the phase transition is first order as is assumed in Fig. 3, then the chiral transition evolves into the T, $e^{-1/m}$ plane. This line might either terminate somewhere in this plane, as is assumed in Fig. 3a, or it might join on to the upper axis. The most elegant possibility is that the line emanating from the deconfinement transition joins the chiral transition. In this sense, the confinement-deconfinement phase transition and the chiral phase transition would be identified. There are of course many more phase diagrams than are illustrated in Fig. 3, and the actual diagram realized in nature is not yet known, although recent Monte-Carlo simulations favor the possibility that the chiral and deconfinement phase transitions are joined together. It is also important to recognize that although quark masses may be varied in a theory, in nature masses are fixed, and the issue of whether or not there are true phase transitions in QCD is whether or not for some fixed values of masses, the first order phase lines are crossed while varying the temperature T.

In spite of all the theoretical understanding of the properties of hadronic matter which has developed in the last few years, almost no experimental data exists which might probe its properties. Such data might provide a striking confirmation of ideas about confinement and chiral symmetry breaking. The production and study of new forms of matter is also of intrinsic scientific interest, and may enlighten conjectures about neutron star structure and the early universe.

There is a good possibility that new forms of matter might be produced and studied in the collisions of ultra-relativistic nuclei. Such a collision is visualized in Fig. 1. In Fig. 1a the two nuclei approach one another in the centerof mass frame. The two nuclei are Lorentz contracted to a thickness $\Delta x \sim 1$ Fermi. This limiting thickness, $\Delta x >> (Mc^2/E_{CM})R$, where R is the nuclear radius, M is the nucleon mass, and E_{CM} is the center of mass energy per nucleon, arises because the nucleons couple to low longitudinal momentum degrees of freedom, $\Delta p^{\shortparallel} \sim 200$ Mev/c, and the Heisenberg uncertainty principle guarantees fuzziness on a scale $\Delta x \sim h/\Delta p^{\shortparallel}$. The longitudinal momentum is the component of momentum along the collision axis. These low longitudinal momentum degrees of freedom are the mesons which are produced with small momentum in the collisions of hadrons with high center of mass energy. The nucleon degrees of freedom, that is the portion of the nuclear wavefunction which carries baryon number, is Lorentz contracted to a much smaller size $\Delta x \sim (Mc^2/R_{CM})R$.

As the two nuclei pass through one another, Fig. 1b, the low momentum meson degrees of freedom interact and a hot hadronic matter distribution forms. Meson degrees of freedom with higher longitudinal momentum take longer to interact, and Lorentz time dilation slows matter formation. As time proceeds, this higher longitudinal momentum matter forms at

$$ct \sim \hbar p''/m^2c^2 \sim x'', \tag{4}$$

where we assume mesons have large p'' and travel with velocities close to that of light, and m is a typical hadronic mass scale, $m \sim 200$ Mev. In this formula, the particles are assumed to travel as free massive particles before they begin to interact significantly, so that $x = ct$. The distance x is given as the Lorentz time dilation factor p''/mc, valid for ultra-relativistic particles, times the Compton wavelength of the forming particle, $\hbar/mc$. The particle motion is assumed to be primarilly along the collision axis, which is true for particles produced in very high energy collisions.

In the rest frame of the newly formed matter, we are implicitly assuming there is an intrinsic interaction time which is the time it takes light to travel the Compton wavelength of the produced particle. This time is essentially the same for all the produced particles since they are almost all pions and have the same mass. Since the matter does nothing until there is time for interactions, we say that the matter has not 'formed' until the interactions begin. As time evolves, more and more matter forms at increasingly large distances from the collision point. The formed matter literally chases after the nuclei, and eventually catches up to them. The nuclei are the last particles to interact since it is expected that the nucleons in the nuclei will form the fastest fragments produced in the center of of mass frame. The time dilation of nucleon interactions within the nuclei is the largest of all produced particles. The matter formed in these nuclear fragmentation regions is baryon rich whereas the matter in the central region was primarily composed of mesons.

A consequence of the inside-outside cascade model is a correlation between space-time and longitudinal momentum. This correlation may provide a powerful analytical framework for ultra-relativistic nuclear collisions.

This picture of hadronic interactions describes hadron-hadron, hadron-nucleus, and nucleus-nucleus collisions. Why are nucleus-nucleus collisions better suited to producing a quark-gluon plasma than are ordinary hadron collisions? The most obvious reason is that in nucleus-nucleus collisions, the average number of produced particles is higher than for hadron-hadron collisions, and higher energy densities result on the average. Another is that nuclei are of large spatial extent, and in order to produce a well thermalized matter distribution, the particles produced in the collisions must rescatter. The large spatial size enables the produced particles to rescatter several times before being emitted from the collision region. Yet another reason is that with the large average number of particles produced in nucleus-nucleus collisions, the statistical fluctuations in the distributions of particles produced in individual events are smaller, and collective phenomenon may be easier to observe. The existence of these collective effects may be crucial for arguing that a quark-gluon plasma has been produced in the collision.

On the other hand, fluctuations in hadron-hadron collisions might also produce a quark-gluon plasma. In very high multiplicity hadron-hadron collisions, very high energy density regions may be produced. The energy density might be so high that even at late times when the system is in a relatively large spatial volume, there may be sufficiently large energy densities to maintain a plasma.

If the space-time volume is sufficiently large that hydrodynamics may be used to describe a collision, the description greatly simplifies. Hydrodynamic equations require knowledge of the state of the matter at some time, and the equation of state of the matter. Monte-Carlo simulations may soon provide good equations of state for hadronic matter, and the equation of state is already well understood with semi-quantitative models. The initial conditions are semi-quantitatively predicted by models of the collision dynamics, and the final state of the matter may be experimentally measured. The evolution of the matter at intermediate times appears to be computable, a situation which is much advanced relative to our knowledge of hadron-hadron collisions.

Assuming an initial formation time for the hadronic matter produced in a nucleus-nucleus collision, and knowing the final state distribution of the matter, the hydrodynamic equations may be integrated backwards in time to determine the initial energy density.(2) There are several experimental measurements of the final state hadron distribution in nucleus-nucleus collisions provided by the JACEE cosmic ray experiment.(5) Extrapolating their results to the head on collisions of Uranium, the energy density is

$$E/V \sim \frac{10 \text{ Gev/Fm}^3}{\tau_0{}^{4/3}} \tag{5}$$

where τ_0 is the intrinsic hadronic interaction time. This result only depends on the isentropic nature of the hydrodynamic expansion, and not on the details of the equation of state. Theoretical estimates of the interaction time τ_0 vary, but the range of reasonable values is probably $.1 < \tau_0 < 1$ Fermi/c. The corresponding range of energy densities is 10-200 Gev/Fm^3. This is 50-1000 times the energy density of nuclear matter and is probably sufficient to produce a quark-gluon plasma.

There are of course great uncertainties in this prognostication. Is the picture we outlined above really appliciable at the energies appropriate to the JACEE experiments? Are non-equilibrium effects not included in the hydrodynamic analysis important? Are the few results of the

JACEE analysis typical of nucleus-nucleus collisions? These questions are difficult to assess without accelerator experiments.

There is of course a pressing question about the signals of a quark-gluon plasma if produced in a nucleus-nucleus collision. There have been many suggestions.[2] Metastable droplets of quark-gluon matter might produce long lived droplets of dense matter, or might be short lived and detonate explosively. Large scale density fluctuations, like steam bubbles in boiling water, might form as the matter non-explosively cools, or as the plasma is formed from hadronic matter. Leptons and photons might probe the plasma in the hot early stages of its existence. Strange particle production might be copious, and there might be abundant Λ production in the fragmentation region. These signals for plasma production are extremely speculative, and experimental data combined with more quantitative theoretical studies are needed to channel speculation into quantitative science.

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Figure Captions:

Fig. 1: Ultra-relativistic nuclear collisions in the center of mass frame. (a) Approach (b) Passage through one another and production of a hot central region. (c) Nuclear fragmentation and cooling of the central region.

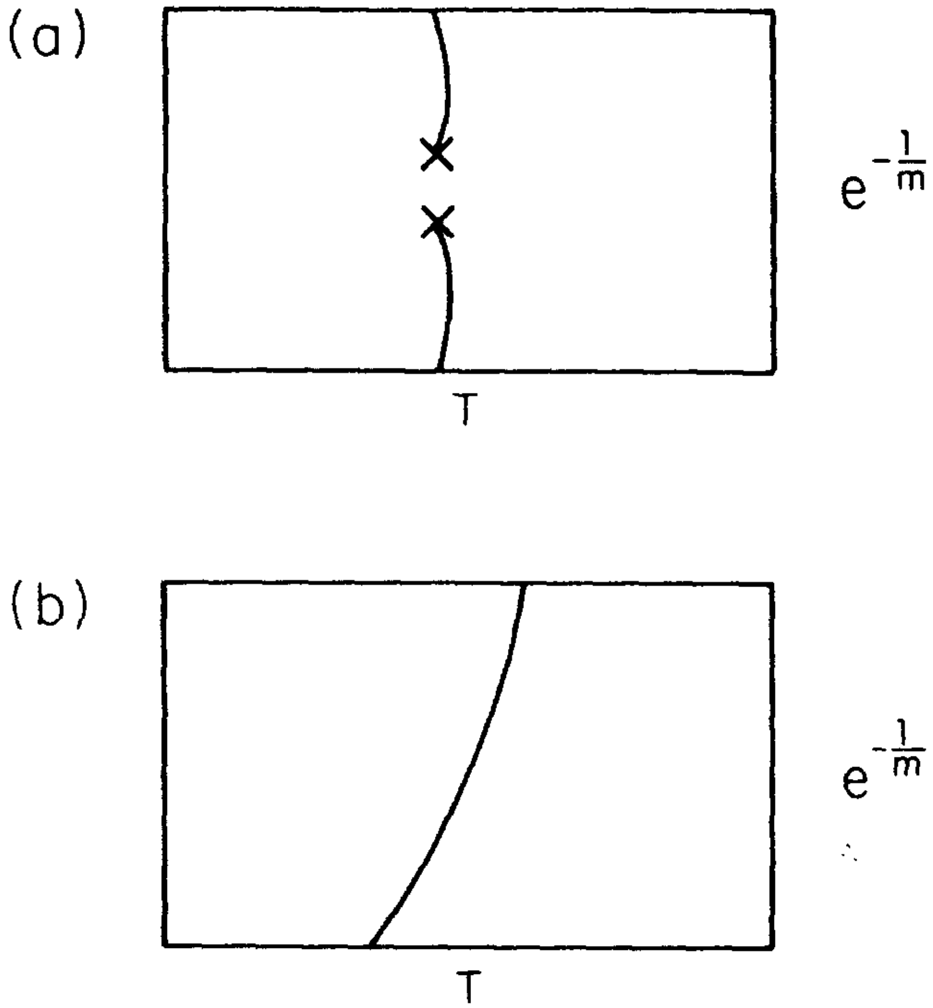
Fig. 2: A representation of results from lattice Monte-Carlo studies. (a) The expectation value of L as a function of temperature. (b) The energy density as a function of temperature.

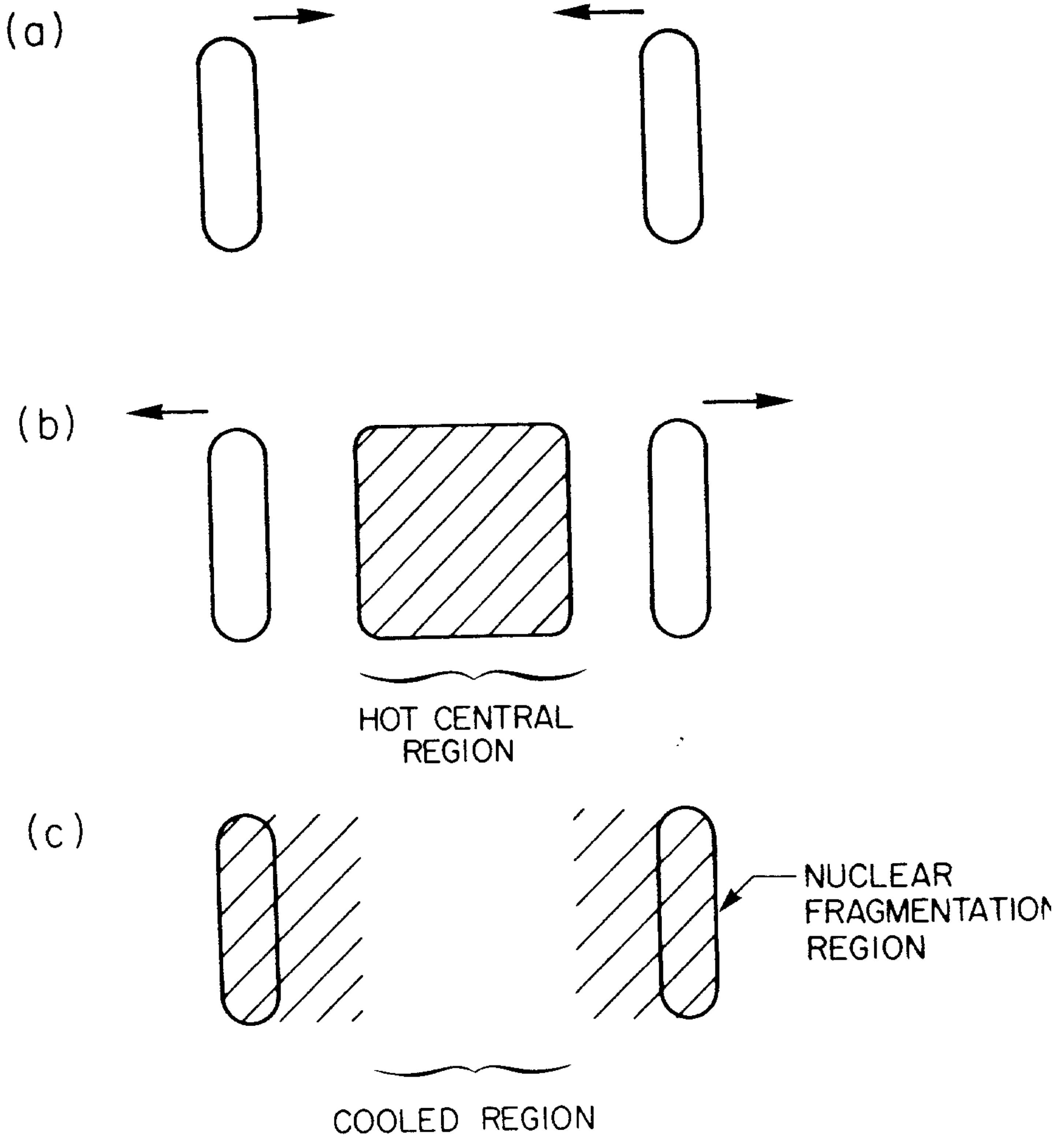
Fig. 3: Hypothetical phase diagrams in the T, $e^{-1/m}$ plane. (a) The chiral deconfined world is not isolated from the chiral symmetry broken, deconfined world. (b) The chiral symmetric, deconfined world is separated from the confined, chiral symmetry broken world.

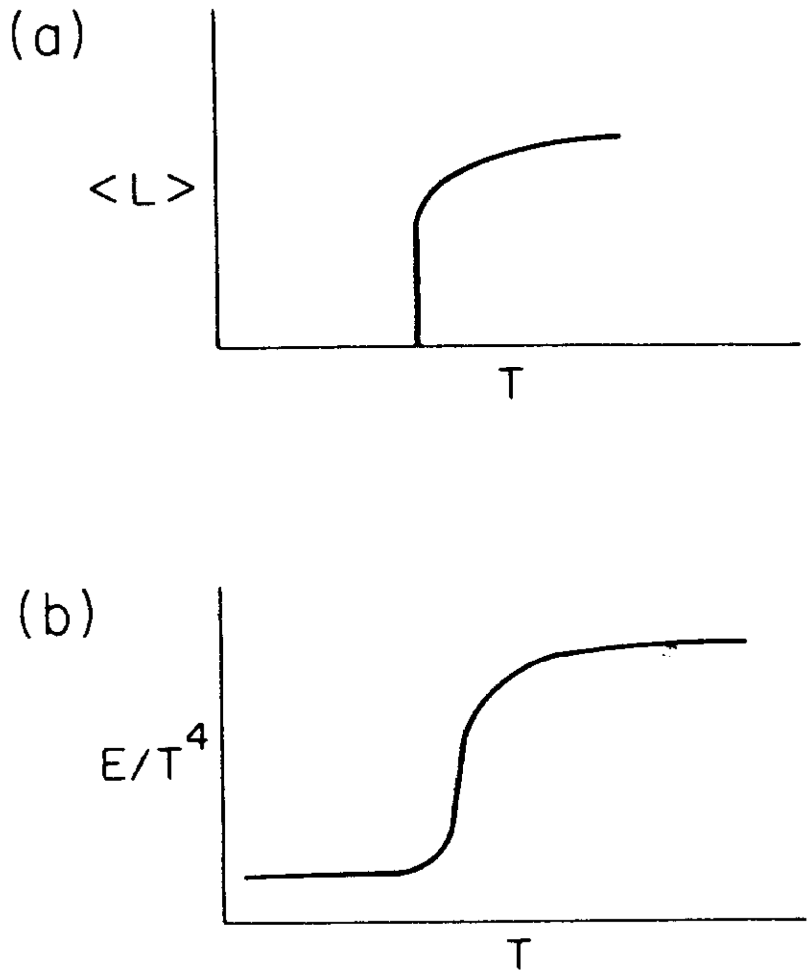
Table 1

Energy Densities and New Degrees of Freedom

ρ (Gev/Fm^3)	Degrees of Freedom
10^{-9}	Electron-positron pairs, photons
1	Leptons, photons, quarks, and gluons
10^{10}	The above plus Z^0, W, and Higgs bosons
10^{62}	The above plus the hypothetical super heavy particles predicted in Grand Unified Theories
10^{78}	The above and the unknown.


(a)
$e^{-\frac{1}{m}}$
T
(b)
$e^{-\frac{1}{m}}$
T


(a)
(b)
HOT CENTRAL
REGION
(c)
NUCLEAR
FRAGMENTATION
REGION
COOLED REGION


(a)
<L>
T
(b)
E/T4
T