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ASPECTS OF QUARK GLUON PLASMA¹

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ABSTRACT

Identifying reliable signatures of the Quark-Gluon Plasma formation in the ultra-relativistic heavy ion collisions will be the major goal of the future experiments at BNL (RHIC) and CERN (LHC). As these experiments are expected to give new dimension to this field of research, it is worth considering to review the main aspects of the Quark-Gluon Plasma (QGP). In this paper we discuss briefly some of the important features of the ultra-relativistic heavy ion collisions and the physics of the QGP. Finally we make a few suggestions based on the most recent developments in the field, with the hope that perhaps questions to the related matters will be answered in the future investigations.

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1. INTRODUCTION

Of the many phase transitions that are studied in particle physics and cosmology the confinement transition of quantum chromodynamics (QCD) is perhaps the least understood theoretically. Though the non-perturbative features of QCD such as confinement have not yet been rigorously established, it is widely believed that at either high temperature or high densities normal hadronic matter undergoes a phase transition to a new unconfined state of matter, the Quark-Gluon Plasma (QGP) formed by the basic constituents of the hadrons i.e; the quarks and gluons (1-3). At low energy densities, quarks and gluons are bound by the strong force into colorless objects, the hadrons (confinement). In addition, the quarks acquire a large effective mass ($m_u \approx m_d = 300 \text{ MeV}, m_s \approx 500 \text{ MeV}$) via interactions between themselves and the surrounding physical vacuum (broken chiral symmetry). When increasing the energy density by increasing the temperature (heating) or the matter density (compressing), a phase transition might occur towards the QGP, the true perturbative vacuum of QCD, where partons are deconfined and chiral symmetry is approximately restored ($m_u \approx m_d \approx 5 \text{ MeV}, m_s \approx 150 \text{ MeV}$) (4). Estimates of the critical temperature and density are $T_c \approx 150 - 200 \text{ MeV}, \rho_c \approx 1 - 3 \text{ GeV} / \text{fm}^3$ (4).

In the recent years, one of the more active fields in high energy physics has been that of identifying reliable signatures of QGP formation in ultra-relativistic heavy ion collisions (URHIC). The proposed signals like strangeness enhancement, multiplicity fluctuations, p_T enhancement, J/ψ suppression and dilepton emission etc., have been achieved, but explanations other than QGP have been given (5). At present new experiments are being planned at BNL (RHIC) and CERN (LHC) because URHIC are still believed to provide a way of producing the QGP related conditions in the laboratory (6,7). It is still an interesting subject to study whether the QGP phase transition does take place, and if so what is the order of transition. Several are the signatures of the QGP formation which have not been completely rejected, and new theoretical developments are underway. It therefore requires to review the whole situation with respect to the QGP physics and the formation signals to know more about the QCD phase transition before the start of new experiments.

There are several recent reviews of the QGP and URHIC already available in the literature (4, 8-13), which place emphasis on the experimental signals and some on the reaction dynamics. Most recent reviews of models (14,15) deal with Dual Parton Model and its MC code VENUS. We shall therefore take the supplementary approach here, and discuss the underlying ideas of QGP Physics and URHIC, with concentration on the recent developments in the field to suggest what possibly should be the line of action for the future experimental investigations. It is difficult to review the whole QGP literature within the scope of a single article, but we shall touch upon as many issues as we can without going into the theoretical details, making the best of what has been published upto the middle of 1995, and see how far we may go along the road towards the QGP in the future Brookhaven (RHIC) and CERN (LHC) experiments.

2. THE QGP AND URHIC

Ultra-relativistic heavy ion collisions have given us the opportunity to search for the QGP formation in the laboratory (16). There are two basic questions one may ask with the above statement in view: (a) why the heavy ion collisions are significantly required to look for the QGP phase transition, and (b) why is it important to study the QGP phase transition in laboratory experiments? We shall answer these questions very briefly in this section.

2.1 Why Heavy Ion Collisions

The physical environment in which the critical conditions of temperature T_c and density ρ_c can be achieved are known to be: the early universe (17), the core of neutron stars (17), and the heavy ion collisions (18,19). It can immediately be understood that the best chance for actually studying the properties of QGP is in the laboratory by colliding the heavy ions (18,19). In the heavy ion collisions one has an access to soft events with unusually large multiplicities where the energy is distributed as evenly as possible among a very large number of secondaries, to probe hadronic matter at high density near thermodynamic equilibrium (Strong Interaction Thermodynamics).

In a single nucleon-nucleon collision at AGS or CERN energies, a strange baryon for example is created approximately 15% of the time and the probability of creating several strange baryons in a single nucleon-nucleon collision is therefore substantially smaller (20). In heavy ion collisions, on the other hand, several nucleon-nucleon collisions take place simultaneously and a single Au + Au collision could result in the production of 20 Λ particles in close proximity to each other. A large number of particles produced in this manner in a finite volume of the collision suggests the release of a fairly large energy density during the collision which gives a possibility of the existence of the QGP.

The QGP phase lies between the Early Universe and the Neutron Star regimes on the phase diagram; which we shall discuss it in section 4. Heavy ion collisions at various beam energies can probe the entire region in between these two extremes. Collisions between ultra-relativistic heavy ion collisions thus offer the best chance to create and observe the QGP in the laboratory. For a given accelerator energy and thus centre-of-momentum (cm) energy, nuclei that are as heavy as possible produce the highest energy density. To enhance the possibility of observing the QGP, we want high energy density over a sufficiently large volume, with dimensions larger than the scattering length of the constituents, so that equilibrium thermodynamics may be applicable. Once created, the system must exist long enough to equilibrate so that a phase in the thermodynamic sense can be established. These considerations led to the experimental programs at Brookhaven (AGS) and CERN (SPS). At each of these facilities it was expected to create a QGP in a baryon-rich environment. The quest for a QGP in a region with no (or a very low) net baryon density, but significantly higher energy density, led to the decision to build the Relativistic Heavy Ion Collider (RHIC) at Brookhaven which is under construction and the Large Hadron Collider (LHC) at CERN to be constructed.

2.2 The Study of QGP

Quantum chromodynamics (QCD) predicts the existence of QGP, a new phase of strongly interacting system, at high energy density. The discovery and proper understanding of QGP is thus of much significance for QCD since it fortells the long-range behaviour where theory is still poorly understood. This phase transition

has not only intrinsic importance but also its finding in the laboratory could enlighten the development of the collision process in space-time.

The early universe presumably underwent such phase transition $10^{-10} - 0.5$ second after the Big Bang (21). Attempts to create the QGP in the laboratory by a “little bang” will therefore shed some light on the “Big Bang”, and hence on the history of early universe. Phase transition at high temperature is also useful in the cosmological studies. Critical phenomena that can occur, for example long-range density fluctuations (as in condensing water), might have a bearing on important aspects of cosmology, such as nucleosynthesis, dark matter and the large-scale structure of the universe (22). There are speculations that a strong first order (or discontinuous) phase transition would result in relics of gravity waves, or a new form of matter or alter nuclear abundances, or baryon density inhomogeneities in the cosmic fluid (23). In astrophysics the dynamics of supernovae explosions depends on the compressibility and therefore the equation of state of the nuclear matter (24). It is even speculated that the cores of neutron stars may consist of cold QGP. The QGP phase of matter thus also provides key information about the inner cores of neutron stars having larger nuclear density, $\rho/\rho_{\text{nuc}} \approx 10$ (25).

We now discuss these aspects in some more detail

2.2.1 The Study of Cosmology

May be the most compelling reason why we should attempt to study the Quark-Gluon Plasma transition in laboratory experiments is that it must have occurred in the early universe (26). The history of the thermal evolution of our universe is depicted in Fig.1. The relation between temperature T and cosmic time t is approximately given by (27):

$$T_{\text{MeV}} \approx (5.75 N_f(T))^{1/4} t_{\text{sec}}^{-1/2},$$

where $N_f(T)$ describes the number of particle degrees of freedom that act as effectively massless modes at a given temperature, and the subscripts indicate the units in which T and t are measured. From the present temperature of the cosmic background radiation (2.7 K) we extrapolate back to temperature of about

2×10^{12} K $\approx$ 200 MeV at about 20 μ s after the “Big Bang”. This is the temperature above which hadrons dissolve and their constituent quanta, quarks and gluons, are liberated.

Tracing the history of our universe backward, this is only the first phase transition involving fundamental quantum fields that we encounter. Most likely, more transformations of a similar nature have occurred at even earlier times. If our current ideas of the origin of electroweak symmetry breaking are correct, a phase transition in the Higgs vacuum took place at $t \approx 10^{-11}$ s, when the temperature was around 250 GeV. A similar phase transition at much earlier times ($t \approx 10^{-35}$ s) associated with the “grand” unification of electroweak and strong interactions may have led to exponential inflation of cosmic scales, due to the anti-gravitational pressure exerted by an unstable vacuum state (28,29).

Why should we care about events occurring at such unimaginably short moments after the creation of our universe. The reason is that some of the unsolved problems of cosmology are, most likely, associated with events in the era before 10^{-5} s:

- The observed *baryon number asymmetry* in the universe is probably due to baryogenesis by topologically nontrivial field configurations during the electroweak phase transition.
- The nature of *dark matter* may be associated with properties of the Higgs or the QCD vacuum.
- The near *isotropy of the background radiation*, and the deviations from the homogeneity underlying the observed *large scale structure* of the universe were according to current thinking, determined during the phase transition causing inflation.

Understanding the dynamical nature of phase transitions in elementary quantum field theories is thus essential to further progress toward the solution of these puzzles. Since the QCD phase transition is the only one of these that appears accessible to laboratory experiments, we must study it with higher priority.

2.2.2 The Quark Stars

In addition to events in the very early universe, quark matter may also play a role in the internal structure of collapsed stars. At the high densities reached in the core of neutron stars nucleons may well dissolve into their constituents, forming baryon-rich cold quark matter. This would not lead to greater stability of neutron stars, quite to the contrary: since quark matter would allow for a higher central density of a star at a given total star mass, its formation would actually facilitate collapse to a black hole. Neutron stars with a quark core have a lower value of their upper mass limit, probably somewhere between 1.5 and 2 solar masses (30). Stars with a quark core would be more compact and hence could sustain higher rotation rates (31). This observation would be of practical interest, if pulsars with periods in the sub-millisecond range are eventually detected.

The problem with quantitative predictions is that dense baryon-rich nuclear matter is expected to contain a large strangeness fraction, because the inclusion of strange quarks can lower the Fermi energy. This holds true for baryonic matter as well as quark matter. Unfortunately the equation of state of baryonic matter containing hyperons is poorly known. The scalar coupling strength of Λ -hyperons, which is not well determined experimentally has a large influence on the central density of hyperon-rich neutron stars. This uncertainty can easily mask the formation of a quark matter core (30). Better understanding of the equation of state of hyperon interactions in dense baryon-rich matter, which can only come from high-energy nuclear collisions, is therefore essential for further progress towards the understanding of neutron star structure.

3. THE MICROSCOPIC MODELS

One of the main objectives of the new experiments at Brookhaven and CERN will be to identify reliable signatures of the QGP formation. Thus a thorough understanding of soft physics appear to be mandatory if plasma signatures are to be separated from "conventional background". But, there is not a theory available to describe the background necessary to amplify the phase transition signals because we do not know how to apply QCD to soft processes (small momentum transfer), which are the bulk of the events taking place in heavy ion collisions. Luckily, however, the QCD based phenomenological models have been introduced

by extrapolating the ideas developed for the hadron-hadron and hadron-nucleus interactions correctly describing the soft nucleus-nucleus collisions, the best way to analyse the data. The quark-parton model (32), the quark model (33), the additive quark model (34), the constituent quark model (33), the multichain model (35), and the wounded nucleon model (36), are a few examples to quote that look at the dynamics of quark matter. The latest and the more popular, however, are the Lund Model (37), and the Dual Parton Model (38,39). There are several versions based on these models. For Dual Parton Model (DPM) they are IRIS (40), VENUS (41), QGSM (42), DTUJET (43) and Dual Multichain Monte Carlo (44). Closer to DPM is PYTHIA (45) and HIJING (46). The DPM Monte Carlo use different codes for fragmentation functions, but are all applicable to nucleus-nucleus collisions. Lund Model versions are HIJING (46), FRITIOF (47), ATTILA (48), SPM (49), COJET (50), EUROJET (51) and RQMD (52).

3.1 The Basic Ideas

In contrast to the hydrodynamical/thermodynamical models which are based on coherent nucleus collisions: each nucleus interacts as a whole, the microscopic models take into account the incoherent mechanism: there is no collective effect and the collision is a succession of independent nucleon-nucleon collisions. These models are altogether non plasma models, which have string formation and fragmentation as an important ingredient. They attempt to describe the phenomenology of heavy ion collisions based on extrapolation from known region of high energy interactions, and might thus be used to discriminate "new physics" from the background of conventional physics. These models have been created to study what is expected in the ultra-relativistic heavy ion collisions. Any deviation of the experimental data may signal the onset of a new physics.

These models have now been further developed on the basis of QCD inspired mechanism incorporated in Dual Parton Model and Lund Model Fritiof which consist in a string formation followed by a fragmentation of these strings in observed hadrons. Interaction between strings have been modelled as color rope formation in RQMD (53), string fusion in QGSM (54) or quark matter droplet formation in VENUS (55). On the hand, collectivity in flavor production may also emerge in the hadronic stage of nuclear reactions due to RQMD (52) and ARC (56).

3.2 Interaction Mechanism

Three of the theoretical models VENUS (41), RQMD (52) and ARC (56), as well as the Lund Model Fritiof (47) have been so successful that experimentalists nowadays routinely refer to them when they want to know whether their data exhibit something unexpected (57). We therefore will discuss their main interaction mechanism without going into the theoretical details. The reader is referred to the original papers and the recent reviews (4, 14-17) for an extensive account.

3.2.1 Collision geometry

The geometry of the collision is treated in most models in a very similar fashion. The nucleons of the projectile and target nuclei move through each other on straight-line trajectories. Interaction between two nucleons occur whenever they come closer than some minimal distance, e.g. $d < (\sigma_{nn}/\pi)^{1/2}$. The results of the reaction is a longitudinally oriented object called a 'string'. A string can be conceived as a quark-antiquark or quark-diquark pair which stretches an attractive 'color field' or 'flux tube' between them.

3.2.2 String formation

The detailed mechanism of string formation is different in the various models. VENUS, IRIS, and QGSM etc. form strings via color exchange, while FRITIOF, ATILA, and RQMD creat them via momentum exchange (longitudinal excitation). This is illustrated in Fig.2. In the case of color exchange, a target parton is linked to a projectile diquark and vice versa. Neither a particle nor momentum has been exchanged, but a color exchange has formed two new strings which have acquired large (longitudinal) masses because of the large momentum difference between the newly combined partons. In the case of longitudinal excitation the strings are stretched between the original partons of the projectile or the target nucleon. The required excitation of the string comes from a longitudinal momentum transfer between the hadrons. Despite the difference in their formation, the resulting strings are quite similar; measurable differences are provided only by the different flavor content of the strings. It has been demonstrated by Werner (58) that the inclusion of anti-quark color exchange and diquark break-up increases the

baryon content in the central region. No transverse momentum transfer is considered at this stage of the reaction. This is in accord with the 'soft' character (only a few hundred MeV trasverse momentum is carried by the final hadrons) of the hadron-hadron interactions the models are made for.

3.2.3 String fragmentation

Transverse momentum and mass are assigned to string pieces or hadrons at string fragmentation. Again there are different prescriptions on how a string breaks up into smaller string pieces or hadrons. VENUS employs the Atru-Mennessier formalism (59,60) for string fragmentation; IRIS, FRITIOF, ATILA and RQMD use the LUND formalism (61); QGSM, and earlier versions of VENUS use the Field-Feynman model (62) for fragmentation. The differences between the fragmentation models are mainly of a technical nature, e.g. how to determine the location of the breakpoints on the string surface.

3.2.4 Rescattering and the concept of a formation time

Owing to the similarities of the employed formalisms, the calculated observables from the various models are very similar (63). Considerable differences emerge if rescattering of secondary particles is included in the models. Two aspects of rescattering are particularly important: the excitation of the target spectator and the modification of particle yields. At present, VENUS, QGSM, and, RQMD have implemented reinteractions of produced particles. The interaction of produced particles is closely tied to the concept of 'formation time' (64,65) of secondary particles. The formation time cannot be calculated in the QCD because hadronization is a soft process, belonging to the non-perturbative sector. QGSM treat secondary-particle production as a black box as far as the formation time is concerned: the hadrons are created with a certain spectral distribution at a space-time point (x,t) (= the collision point) and, from then on, propagate freely to the point $(x',t+\tau)$, where they become observables (on small-shell states). During this formation time τ , the preformed hadron does not reinteract with the surrounding hadronic matter. The proper formation time τ_0 is given by

$$\tau = E / m\tau_0$$

In principle, τ_0 is now a parameter which has to be determined by comparison with experimental data. Within the string picture, Bialas and Gyulassy (66) have introduced a constituent and a yoyo-formation time, τ_c and τ_y , respectively. τ_c is the time when the string breaks, i.e.; new constituents are created; τ_y is the time when the 'yoyo state' is formed. Both VENUS and RQMD contain interactions of constituent (di)quarks with hadrons which occur for times $\tau_c < t < \tau_y$. Leading (di)quarks may also interact for times $t < \tau_c$. Thereafter one has hadron-hadron interactions.

For simplicity, VENUS assumes the cross-section for (di)quark-hadron and hadron-hadron interactions to be identical. They are parametrized in terms of a unique interaction radius r_0 . Again, r_0 can, in principle, be determined from comparison with experimental data.

In the RQMD model the interaction distance is, where known, defined by the experimentally measured cross-sections of each individual hadron-hadron collision. Otherwise, hadronic interaction models, such as the additive quark model, boson exchange, the Breit-Wigner formalism or Regge phenomenology, are employed. The result of an interaction is either strings (depending on the energy), or resonances as they are defined in experimental tables with their measured lifetimes, branching ratios and masses.

Thus, most models are similar in the respect that they parametrize the 'strength' of rescattering with a single parameter, which could be determined from comparison with experimental data. An exception is the RQMD model which includes, as far as possible, measured cross-sections and decay probabilities. VENUS, QGSM and RQMD include rescattering of all particles.

Owing to the large number of particles ($> 40\,000$) which are predicted for heavy-ion collisions at LHC energies ($\sqrt{s}=6.3$ TeV), rescattering as sequence of two-body scatterings becomes undefined. A new 'cluster-ansatz' has recently been proposed by Aichelin and Werner (67), where overlapping hadrons are treated as clusters with a particular excitation energy.

4. PHYSICS OF HEAVY ION COLLISIONS

Collisions between relativistic nuclei are studied in order to explore nuclear matter in the entirely uncharted regions of high energy and baryon density. One is interested in the space-time development of the hadronic and partonic interactions in nuclear volumes under such conditions.

In the context of the Standard Model, the study of this phase diagram of strongly interacting matter (Fig. 3) is not only of interest in exploring and testing QCD on its natural scale, i.e. in the non-perturbative sector, but it might also shed light on such fundamental questions as the nature of confinement itself and on the process of spontaneous symmetry breaking, which is made responsible for the origin of the 'effective' quark masses (the pion being the assorted Goldstone boson).

In order for heavy-ion collisions to fulfil these promises, a number of necessary preconditions have to be met and verified by the results.

- (i) In order to use macroscopic variables, the system created has to be 'big', i.e.; its dimensions ought to be much larger than the typical scale of strong interactions ($\gg 1$ fm) and it should consist of 'many' particles ($\gg 1$).
- (ii) In order to use the language of thermodynamics, the system has to be in (or near) equilibrium, i.e. its life time has to be larger than typical relaxation time ($\tau \gg 1$ fm c^{-1}). Thermodynamical concepts are frequently applied in the description of heavy-ion collisions, but their applicability has to be verified and checked with the microscopical, non-equilibrium calculations. Equilibrium can be reached and maintained throughout the expansion only in the sufficiently interacting system, hence the number of collisions per particle has to be larger than one. Rescattering of the produced particles (hadron/quarks) amongst themselves and with the surrounding nuclear matter is therefore of vital importance and not a trivial effect or even a nuisance. In fact only a few (typically > 3) collisions are required, for example, to equilibrate momentum distributions in low-energy nucleus-nucleus collisions (68).
- (iii) It has to be ascertained that the energy densities of the order of $1-3$ GeV fm^{-3} , predicted for QGP by QCD, can indeed be reached in heavy ion collisions.
- (iv) Because the created system is not static but, rather, rapidly evolving, experimental observables will, in general, correspond to an integral over the

complete space-time history of the reaction until freeze-out, and disentangling the various contributions to a signal from the different phases shown in Fig. 3 presents a formidable challenge indeed. Furthermore, a system evolving in equilibrium by definition erases its memory of preceding stages. As a Big Bang cosmology, it is necessary to identify observables that decouple at different times from the expansion and are more sensitive to the early and hot stages of the matter.

4.1 Space-time Evolution of Heavy Ion Collisions

Nuclei are extended objects and therefore their geometry plays an important role in heavy-ion collisions. It has been shown (69) that since inelastic part of proton-proton cross-section only varies with a few percent in the 1 GeV to 1 TeV range, the geometry in ultra-relativistic heavy ion collisions is basically independent of the incident energy. It has actually been found that both the slow target associated particles (70), the projectile fragments (71), as well as the particles produced in the fragmentation region (72) show energy independent features. Thus nuclear geometry is said to be energy independent over a large range of energy, and a clean-cut particle spectator picture may well be used as a guiding hypothesis for ultra-relativistic heavy ion collisions as assumed in most of the models (73). Projectile mass independence has also been shown in the emulsion experiments (74). The WA80 collaboration (73, 75-76), using a simple geometrical approach, calculated average transverse momentum $\langle p_t \rangle$ per participating nucleon in different type of interactions. The results from these analyses clearly indicate that, for a given geometry, the number comes out essentially the same independent of the target mass, projectile mass, and centrality.

Fig. 4 shows the sketch of a reaction between asymmetric nuclei A and B, the impact parameter b separates the nucleons into *participants* with primary nucleon-nucleon collisions, and *spectators*, which proceed with little perturbation along the original direction. The use of this straight-line is justified by the fact that at very high energies the size of nucleons is larger than their Compton wavelength and the nuclear radius is larger than the interaction length of ≈ 1.8 fm. Cross-sections, number of participants, and related kinematic quantities are usually calculated from simple geometrical considerations.

The space-time evolution of central ($b = 0$) collision at very high energies (RHIC & LHC) is shown schematically in Fig. 5. In the first moments of the reaction, hard scattering processes on the parton level may occur with a (small) probability given by structure functions and perturbative QCD cross-sections. In addition, soft nucleon-nucleon collisions (large cross-section) between the two highly Lorentz-contracted nuclei (with limiting thickness $r \approx 1$ fm) (77) redistribute a fraction of the original beam energy into other degrees of freedom.

After a short time usually taken to be of the order of ~ 1 fm c^{-1} ('formation time') (78), partons materialize out of the highly excited QCD field. Thermal equilibrium may now be approached via individual parton-parton or, equivalently, string-string interactions. Calculations of the mean free path of the quarks in QCD matter give a value of $\lambda \approx 0.5$ fm at energy densities $\epsilon = 2$ GeV fm $^{-3}$ (79), thus indicating that equilibrium may indeed be reached in collisions of heavy nuclei, where the transverse radii, and hence the initial dimensions, are clearly larger than λ .

However, the system expands rapidly, mainly along the longitudinal direction, and therefore lowers its temperature and reaches the critical temperature region T_c after $\tau \approx 3-5$ fm (5). Potentially the matter now spends a long time in the mixed phase ($\tau > 10$ fm c^{-1}), in particular if the transition is of first order (79). It has to rearrange the many degrees of freedom (partons) of the QGP into the smaller number available in the hadron phase, with large associated release of latent heat. In the last and hadronic phase ('hadron gas' or 'hadron fluid', $\tau \gg 10$ fm), the still interacting system keeps expanding, possibly in an ordered motion ('flow'). It may expand at very large dimension ($V > 10^4 - 10^5$ fm 3) until 'freeze-out', when interactions cease and the particles stream freely away to be detected in the experiments.

The timescale introduced above is, however, quite uncertain and has, up to date, no experimental justification. The various signals that can be observed by the experiments and associated (to some extent) with the different stages of the evolution will be discussed in the following section.

5. PHYSICS OF THE QUARK-GLUON PLASMA

Although it is generally believed that a phase transition should take place from normal hadronic matter to the unconfined state of quarks and gluons at high energy density, many aspects of this transition are still to be understood theoretically. Here we give a brief survey of (mostly) theoretical aspects of the QGP to create fresh interest in the field. More detailed discussions of selected topics as well as review of experimental results can be found in other recent articles (8-13).

5.1 The Equation of State of Hadronic Matter

There are two approaches that have been widely followed to the problem of the equation of state of strongly interacting matter:

- (a) Consider color-singlet hadrons, i.e. baryons and mesons, only and see how far one gets with taking into account their known interactions and excitation spectra. This approach was pioneered by R. Hagedorn, and later studied by Hagedorn and Rafelski(81), Walecka(82), Gasser and Leutwyler(83), and many others.
- (b) Consider the fundamental constituents of hadrons, i.e. quarks and gluons, and study the equation of state on the basis of quantum chromodynamics. This approach, anticipated in pre-QCD days by P. Carruthers(84), was pioneered by Collins and Perry(17), Baym and Chin(85), McLerran et al(86) and by Shuryak(1).

Here we will start by examining the approach based on hadrons and their interactions, because its results are important in their own right. Any quark-gluon plasma formed in a nuclear reaction will eventually hadronize and evolve through a hadron-dominated break-up phase.

5.1.1 Hot and dense hadronic matter

How far do we get in explaining the equation of state of hadronic matter by considering hadrons alone? Experimental observations of baryon and meson resonances indicate that the mass density of hadronic states grows exponentially,

$$\rho(m) \approx m^a \exp(m/m_0) \quad (1)$$

as originally postulated by Hagedorn (87). It is now understood that an exponential mass spectrum is the natural consequence of quark confinement, and it is found in even the simplest hadron model that incorporate the confinement concept, such as the string model (88) or the MIT-bag model (89). A numerical simulation of the spectrum of excited string modes agrees very nicely with the observed hadron spectrum even at low energies (90).

Now consider highly excited hadronic matter, characterized by a temperature T . The energy density of excited states is then given by;

$$\begin{aligned} n(E) &\equiv \exp(-E/T) \int_0^E dm \rho(m) E \sqrt{E^2 - m^2} \\ &\equiv E^b \exp \left[-E \left(\frac{1}{T} - \frac{1}{m_0} \right) \right] \end{aligned} \quad (2)$$

Obviously, $n(E)$ is only integrable as long as the factor in the exponent remains positive, i.e. for $T < T_c \equiv m_0$. From experimental data we know that $m_0 \approx 200$ MeV, i.e. the temperature of a (noninteracting) hadronic resonance gas cannot exceed the Hagedorn temperature $T_c \approx 200$ MeV. Cabibbo and Parisi pointed out in 1975 that the distribution remains integrable at $T = T_c$ when $b < -2$, and then corresponds to a finite energy density $\epsilon = \int E dE n(E)$ at the critical point (91).

This is reminiscent of the singular behavior of thermodynamical quantities at a first order phase transition, and they speculated that it signals the transition to a color deconfined quark-gluon plasma.

When Hagedorn and Rafelski (92) studied the effect of the finite internal size of hadrons by utilizing the finite excluded volume approximation, they found that the equation of state developed a singularity at finite energy density at a slightly lower critical temperature $T_c < m_0$, independent of other details of hadronic mass spectrum. We conclude from all this that hadronic matter as composed of

individual color-singlet hadrons ceases to exist at temperatures exceeding $T_c \approx 150\text{-}200 \text{ MeV}$.

5.2 Quark-Gluon Plasma

The theory of the equation of state of quark matter is conceptually much simpler, because it is directly based on the fundamental QCD Lagrangian

$$L_{QCD} = -\frac{1}{4} \sum_a F_{\mu\nu}^a F^{\mu\nu a} + \sum_{f=1}^{N_f} \bar{\psi} (i\gamma^\mu \partial_\mu - g\gamma^\mu A_\mu^a \frac{\lambda^a}{2} - m_f) \psi \quad (3)$$

where the subscript f denotes the various quark flavors $u, d, s, c, \dots$ etc. and the nonlinear glue field strength is given by

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + gf_{abc} A_\mu^b A_\nu^c \quad (4)$$

QCD predicts a weakening of the quark-quark interaction at short distances (or high momenta Q^2), because the one-loop series for the gluon propagator yields a running coupling constant

$$g^2(Q^2) = \frac{16\pi^2}{(11 - \frac{2}{3}N_f)\ln(Q^2/\Lambda^2)} \xrightarrow{Q^2 \rightarrow \infty} 0 \quad (5)$$

The QCD scale parameter is now quite well determined (93) to be $\Lambda \approx 150 \text{ MeV}$. The vanishing of the QCD coupling constant at short distances, called “asymptotic freedom”, has often been taken to imply that interactions among quarks and gluons are negligible in the limit of high temperature or high density. As we shall see below that is not the whole truth, because long-wavelength modes continue to have an important influence on the properties of the quark-gluon plasma

5.2.1 Equation of state of quark matter

So let us first suppose that interactions among quarks and gluons are negligible at high energy density and see what we get. At temperature T and quark chemical potential μ (one third of the baryochemical potential $\mu_B = 3\mu$), the energy density of free gluons, quarks and antiquarks is (a detailed derivation can be found in ref. 94):

$$\epsilon_g = 16 \frac{\pi^2}{30} T^4 \quad (6a)$$

$$\epsilon_g + \epsilon_{\bar{q}} = 6N_f \left(\frac{7\pi^2}{120} T^4 + \frac{1}{4} \mu^2 T^2 + \frac{1}{8\pi^2} \mu^4 \right) \quad (6b)$$

Since we have neglected the quark mass, we inserted a factor N_f counting the number of quark flavors active at a given temperature (essentially those with $m_f \leq T$). Clearly $N_f \geq 2$ at all relevant temperatures, $N_f = 3$ in the range $200 \text{ MeV} \ll T < 1 \text{ GeV}$ where strange quarks can be considered as light, as well. The other thermodynamical quantities of interest, i.e. pressure P , entropy density s , and baryon number density ρ_B , are easily obtained from eqs. (6)

$$P = \frac{1}{3} \epsilon, \quad s = \left(\frac{\partial P}{\partial T} \right)_\mu, \quad \rho_B = \frac{1}{3} \left(\frac{\partial P}{\partial \mu} \right)_T \quad (7)$$

Note that the simple relation between ϵ and P holds only as long as all particles are considered massless. The general relation is

$$\epsilon = T \left(\frac{\partial P}{\partial T} \right)_\mu + \mu \left(\frac{\partial P}{\partial \mu} \right)_T - P \equiv Ts + \mu_B \rho_B - P \quad (8)$$

In order to find the location of the phase transition toward hadronic matter we have to incorporate the breaking of scale invariance provided by QCD interactions. The simplest way of achieving this is by invoking the argument associated with the MIT-bag model: Free quarks and gluons can only propagate where the complex structure of the real QCD vacuum has been destroyed. The vacuum realignment costs a certain amount of energy per unit volume, expressed by the MIT-bag constant $\epsilon_0 = B \approx (150\text{-}200 \text{ MeV})^4$ (4,26), and the “wrong” vacuum is endowed with a negative pressure $P_0 = -B$. The relation $\epsilon_0 = -P_0$ is mandated by the Lorentz invariance of the vacuum state. The negative sign is easily understood as signal of the instability of the wrong vacuum state which collapses if not supported by the pressure provided by free partons propagating in the volume filled with it. The equation of the state of quark-gluon plasma then takes the simple form

$$\epsilon = \epsilon_g(T, \mu) + \epsilon_q(T, \mu) + \epsilon_{\bar{q}}(T, \mu) + B; \quad (9a)$$

$$P = P_g(T, \mu) + P_q(T, \mu) + P_{\bar{q}}(T, \mu) - B \quad (9b)$$

A lower limit of stability of the plasma state is obtained by setting $P = 0$. This yields the stability line in the T - μ plane shown in Fig. 6. Of course, the plasma phase becomes unstable against formation of a gas of color-singlet hadrons even earlier, when its pressure equals that of a hadron gas at the same temperature T and chemical potential μ .

In order to explore this aspect further, let us restrict ourselves to the baryon-free case $\mu = 0$, and explore the coexistence between the quark-gluon plasma phase given by Eq. (9a) and hadronic matter, represented by a gas of massless, noninteracting pions

$$\epsilon_\pi = 3 \frac{\pi^2}{30} T^4, \quad P_\pi = \frac{1}{3} \epsilon_\pi \quad (10)$$

ϵ , P , ϵ_π and P_π are shown as functions of T in Fig. 7.

Thermodynamical phase stability requires that the phase with the larger pressure dominates, and phase equilibrium is achieved when $P(T_c) = P_\pi(T_c)$. As Fig. 7b shows, one finds $(T_c) \cong 0.75 B^{1/4} \cong 150 \text{ MeV}$ with $B \cong (200 \text{ MeV})^4 (4.26)$ in this model. Due to the vacuum rearrangement energy B , the energy density between the two phases differs greatly at this point, by the amount $\Delta\epsilon_c \cong 4B \cong 0.8 \text{ GeV/fm}^3$. This simple model obviously predicts a first-order phase transition between the pion gas and quark-gluon plasma with a large latent heat $\Delta\epsilon_c$.

The model is grossly oversimplified because other hadron masses begin to decrease substantially around $T \cong 150 \text{ MeV}$, leading to an increase in the energy density and pressure of the plasma phase. In first-order perturbation theory the modification of the plasma equation of state is (for $N_f = 2$)

$$\begin{aligned} \epsilon = & \left(1 - \frac{15}{4\pi} \alpha_s\right) \frac{8\pi^2}{15} T^4 + \left(1 - \frac{50}{21\pi} \alpha_s\right) \frac{7\pi^2}{10} T^4 + \\ & + \left(1 - \frac{2}{\pi} \alpha_s\right) \frac{3}{\pi^2} \mu^2 \left(\pi^2 T^2 + \frac{1}{2} \mu^2\right) + B \end{aligned} \quad (11)$$

yielding a reduction by about a factor 2 when $\alpha_s = 0.5$.

More reliable predictions concerning this phase transition can at present only be obtained by numerical simulations of the QCD equation of state on a discretized volume of space-time usually referred to as *lattice gauge theory*. In this approach (95) one approximately calculates the partition function for a discretized version of the QCD Lagrangian (Eq. 3) by Monte-Carlo methods. In principle, this technique should accurately describe the quark-gluon plasma as well as the hadronic phase but, in practice, its accuracy especially at low temperature is severely limited by finite size effects and other technical difficulties. Where the numerical results are most reliable, i.e. for the pure gluon theory without dynamical quarks, the calculations predict a sudden jump in the energy density at a certain temperature while the pressure rises more gradually, as shown in Fig. 8.

5.2.2 Properties of the quark-gluon plasma

(a) Thermalization:

For those of us interested in the detection of quark-gluon plasma in nuclear collisions it is imperative to know something about its rate of thermalization. Does it thermalize sufficiently fast, so that a thermodynamical description makes sense? In the picture based on quasi-free quarks and gluons moving through the plasma, thermalization proceeds mainly via two-body collisions, where the color force between the colliding particles is screened, as illustrated in Fig. 9. The technically easiest way of looking at this is to consider the quark (gluon) *damping rate* γ , i.e. the imaginary part of the quark (gluon) self energy. For quarks and gluons of typical thermal momenta ($p \cong T$) one finds (97,98)

$$\gamma_q = -\frac{1}{4p} \text{Im}[\text{Tr}(\Sigma(p))] \cong \frac{2}{3} \alpha_s \left(1 + \ln \frac{1}{\alpha_s}\right) \quad (12a)$$

$$\gamma_g = -\text{Im}[\omega_T(p)] \cong \frac{9}{4} \gamma_q, \quad (12b)$$

where the factor 9/4 reflects the ratio of the quadratic Casimir invariants C_2 of the octet and triplet representations of color-SU(3). Gluons simply have a higher color charge than quarks ($C_2 = 3$ versus $C_2 = \frac{4}{3}$) and therefore scatter more often.

One can argue that a better way to look at a thermalization is to consider the rate of momentum transfer between particles, i.e. weighting the scattering diagrams of Fig. 9. by $\sin^2 \theta$, where θ is the scattering angle. Here one finds (99)

$$\Gamma_q^{(tr)} \cong 2.3\alpha_s^2 T \ln \frac{1}{\alpha_s}; \quad \Gamma_g^{(tr)} = 3\Gamma_q^{(tr)} \quad (13)$$

The *transport rate* $\Gamma^{(tr)}$ is closely related to the shared viscosity of the quark-gluon plasma. Both approaches yield quite similar numbers for values of the strong coupling constant in the range $\alpha_s = 0.2-0.5$. The characteristic equilibrium time, as defined as the inverse of the rates γ_i or Γ_i are

$$\tau_g \cong 1 \text{ fm} / c, \quad \tau_q \cong 3 \text{ fm} / c, \quad (14)$$

i.e. gluons thermalize about three times faster than quarks (100). Initially therefore probably a rather pure glue plasma is formed in heavy ion collisions, which then gradually evolves into a chemically equilibrated quark-gluon plasma. During its hottest phase the QCD plasma is mainly composed of gluons, which are not accessible to electromagnetic probes, such as lepton pairs and photons. The gluon plasma can only be probed by strongly interacting signals, such as charmed quarks or jets.

5.3 Formation of the Quark-Gluon Plasma

5.3.1 Approaches to the thermalization problem

The quark-gluon plasma equilibrates when it is also close to thermal equilibrium. Relativistic heavy ion collisions pose a very different problem: how do the fully coherent proton wave functions of two nuclei in their ground states evolve into locally quasi-thermal distributions of partons as they are characteristic of quark-gluon plasma state? There are mainly two approaches to this problem that have been extensively investigated: (a) QCD string breaking and (b) the partonic cascade.

In the string picture (Fig. 10) developed from models of soft hadron-hadron interactions, one assumes that nuclei passed through each other at collider energies with only a small rapidity loss (on average about one unit), drawing color flux tubes, or strings, between the “wounded” nucleons. If the area density of string is low (not much greater than 1 fm^{-2}) they are supposed to fragment independently by quark pair production on a proper time scale of order $1 \text{ fm}/c$. Most realizations of

this picture are based on the Lund string model (101), e.g. FRITIOF (47), ATILA (48), VENUS (41), QGSM (42) and RQMD (52). When the density of the string grows further, at very high energy and for heavy nuclei, the fusion of elementary flux tubes to “color ropes” has been postulated (102).

Alternatively, a continuum description based on the Schwinger model (1+1)-dimensional QED with heuristic back-reaction — “chromohydrodynamics” — has been invoked to describe the formation of locally equilibrated quark-gluon plasma (103,104). One general aspect of these models is that initially part of the kinetic energy of the colliding nuclei is stored in coherent glue field configurations, which subsequently decay into quark pairs. The flux tube carry no identifiable entropy. The entropy associated with a thermal state is produced in the course of pair creation. In particular, there are no distinction between gluon and quark thermalization.

The parton cascade approach (Fig. 11), whose basic concepts were developed by Boal (105), Hwa and Kajantie (106), and Blaziot and Mueller (107), is founded on the parton picture and renormalization-group improved perturbative QCD. Whereas the string picture runs into conceptual difficulties at very high energy, when the string density becomes large, the parton cascade becomes invalid at lower energies, where most partonic scatterings occur at energies that are too low to be described by perturbative QCD. Nevertheless, the two descriptions may well be just two different formulations of the same physical processes, since there exists a remarkable similarity between the states produced by QCD bremsstrahlung and QCD flux tubes (108).

Let us look at Fig. 5, where we distinguish three regimes in the evolution of an ultra-relativistic heavy ion collision. The physics governing the evolution during the three stages is very different. In this section we will concentrate mainly on the first, i.e. the preequilibrium state, because it is here where much recent progress has been made in our understanding.

5.3.2 Pre-equilibrium parton physics

The complete spectra of phenomena occurring during the approach toward local thermal equilibrium still awaits exploration. Here we will discuss just two aspects:

color screening (109) and charm production (110), because they are dominated by the evolution of the gluon distribution, which is better understood at this moment. The experimentally very interesting contribution to lepton-pair production, which has also been studied (111), is more sensitive to the quark distribution.

(a) *Color screening*(109):

The screening length λ_D of longitudinal color fields is important, because it defines the range over which coherent color-electric fields can extend. If it is larger than the characteristic confinement radius $\Lambda^{-1} \cong 1$ fm, long-range color electric fields arrange themselves as flux tubes. In the one-loop approximation, λ_D is determined by the momentum distribution $f(k)$ of gluons according to:

$$\lambda_D^{-2} = -\frac{3\alpha_s}{\pi^2} \lim_{|q| \rightarrow 0} \int d^3k \frac{|k|}{q \cdot k} q \cdot \Delta_k f(k) \quad (15)$$

A simple result can be obtained when the gluon distribution is exponential in k_T and flat in rapidity y :

$$\lambda_D^{-2} \cong \frac{3\pi\alpha_s}{R_A^2 \tau \langle k_T \rangle} \frac{dN}{dy}, \quad (16)$$

where dN/dy is the rapidity density of gluons. At the earliest time $t_i \cong \langle k_T \rangle^{-1}$, when the scattered gluons become incoherent, the calculations of ref. (109) based on the HIJING code predict

$$\begin{aligned} \lambda_D &\cong 0.4 \text{ fm (RHIC),} \\ \lambda_D &\approx 0.15 \text{ fm (LHC).} \end{aligned} \quad (17)$$

We conclude that the color screening length at LHC energies will be so short, even after the first sequence of parton interactions, that coherent flux tubes cannot develop. The Lund model with its formation of independent strings that break up by formation of quark pairs is simply not applicable under these conditions. At RHIC, on the other hand, the screening length appears to be marginally favoring a parton cascade description, but models with partially fusing flux tubes ("color ropes") may also be able to describe certain aspects of the pre-equilibrium phase of the nuclear collision.

(b) *Charm production*(110):

Charmed quark pairs are predominantly produced in collisions between two gluons. Since the charm production threshold is rather high, about 3 GeV, the rate

of production in the thermalized quark-gluon plasma with temperature $T \cong 300$ MeV is negligible. Conventional wisdom until recently therefore was that most charmed quarks are produced in primary parton interactions (112). In view of the high density of scattered partons with transverse momenta well above 1 GeV at collider energies, one may suspect that there is a sizable contribution to charm production from secondary parton collisions. This is borne out by a calculation of secondary charm production based on the initial scattered gluon distribution predicted by HIJING (110). The additional charmed quarks populate predominantly the central rapidity plateau where initial charm production is reduced by gluon shadowing effects (see Fig. 12). At LHC energies the total yield of secondary charmed quarks may be twice as large as that of primary charmed quarks.

The amount of secondary charmed production is sensitive to the thermalization time of the parton distribution. It depends on the ratio $\langle \sigma_c \rangle / \langle \sigma_{tot} \rangle$, where $\langle \sigma_{tot} \rangle$ is the average total parton-parton cross section that governs thermalization, while $\langle \sigma_c \rangle$ denotes the averaged cross section for charm production. $\langle \sigma_{tot} \rangle^{-1}$ is proportional to the thermalization time τ_{th} . A measurement of the total yield of charmed particles (mostly D-mesons) in the central rapidity region would therefore provide valuable information on the time-scale of thermalization (110).

5.3.3 Plasma evolution and hadronization

Once the quark-gluon plasma has reached the local thermal equilibrium, its further evolution can be described without reference to the parton reactions at the microscopic level. This concept was first quantitatively developed by Bjorken (78). The hydrodynamic equations for an ultra-relativistic plasma with $P = 1/3 \epsilon$ admit a boost-invariant solution describing a longitudinally expanding fireball with constant rapidity density. When transverse expansion effects are taken into account, longitudinal boost invariance is partially destroyed, but the overall picture remains intact. This scenario has been thoroughly studied by a large number of theorists. Because the results are accessible in several fine reviews (113), we will refrain from discussing it here in detail.

An important aspect of the late evolution of the quark-gluon plasma is its hadronization. Mostly it is assumed that the plasma expands and cools until it

reaches the critical temperature $T_c \cong 200$ MeV and then converts into a hadronic gas while maintaining thermal and chemical equilibrium. More detailed descriptions of the dynamics of hadronization have been developed in connection with the problem of strangeness production.

A totally different approach, explored only very recently, consists in following the partonic reactions at a microscopic level, until the parton density has become sufficiently low to permit the formation of individual hadrons (114). A great deal is known about the mechanism of final state hadron production in e^+e^- and NN-scattering but, unfortunately, we do not know whether this knowledge applies to the hadronization of a quark-gluon plasma in bulk. Therefore, the treatment of hadron formation at the end of partonic cascade is fraught with a great deal of uncertainty.

At present, it is hard to judge the merits of either approach on a theoretical basis. Their validity and usefulness simply depends on whether the microscopic processes during hadronization proceed approximately at thermodynamical equilibrium or not. The lesson from nuclear collision processes at much lower energies (below 1 GeV/ μ) has been that both scenarios are possible: Neutron evaporation from highly excited nuclear fragments is well described by thermodynamics, but nuclear multifragmentation, where coulomb forces play a dominant role, requires a more detailed theoretical treatment.

6. QUARK-GLUON PLASMA SIGNATURES

All theory of the quark-gluon plasma would be largely academic if there were no reliable signatures to observe its formation and to study its properties experimentally. It is impossible to present a complete review of quark-gluon plasma signatures here. We will therefore try to capture the essential ideas and current status of the theoretical studies on the most promising quark-gluon plasma signals. Anyone interested in more details is referred to the review of Kajantie and McLerran (16) and to the proceedings of the Strasbourg workshop (115).

6.1 Thermodynamic variables

The basic idea behind this kind of signatures is to measure the equation of state of superdense hadronic matter, i.e. the dependence of energy density ϵ , pressure P , and entropy density s on temperature T and baryochemical potential μ_B . Here one wants to search for a rapid rise in the effective number of degrees of freedom, as expressed by the ratio ϵ/T^4 or s/T^3 , over a small temperature region.

Of course, one requires measurable observables that are related to the variables T, s , or ϵ . It is customary to identify those with the average transverse momentum $\langle p_T \rangle$, and with the rapidity distribution of hadron multiplicity dN/dy , or transverse energy dE_T/dy , respectively (116). One can then, in principle, invert the ϵ - T diagram and plot $\langle p_T \rangle$ as function of dN/dy or dE_T/dy . If there occurs a rapid change in the number of degrees of freedom, one expects an S-shaped curve, as shown in Fig. 13. Models of the space-time dynamics of nuclear collisions need independent confirmation, especially concerning the correctness of their geometrical assumptions. Such a check is provided by identical particle interferometry, e.g. of $\pi\pi$, KK , or NN correlations (117), which yield information on the reaction geometry. By studying the two particle correlation in different directions of phase space, it is possible to obtain measurements of the transverse and longitudinal size, of the life-time, and of flow patterns of the hadronic fireball at the moment where it breaks up into separate hadrons.

6.2 Chiral symmetry restoration

The two most often proposed signatures for a (partial) restoration of chiral symmetry in dense hadronic matter are enhancements in strangeness and antibaryon production. The basic arguments in both cases is the reduction in the threshold of production of strange hadrons (from 700 MeV to 300 MeV) and baryon-antibaryon pairs (from about 2 GeV to almost zero). As Rafelski pointed out over a decade ago, the optimal signal is obtained by considering strange antibaryons, which combine both signatures (118). The enhanced strange quark production in a chirally restored, deconfined quark-gluon plasma (119) leads to chemical equilibrium abundances for all strange hadrons, which would be difficult to understand on the basis of hadronic reactions alone (120).

6.3 Color response function

The basic aim in the detection of a color deconfinement phase transition is to measure changes in the color response function (13):

$$\Pi_{\mu\nu}^{ab}(q^2) = \int d^4x d^4y e^{iq(x-y)} \langle j_\mu^a(x) j_\nu^b(y) \rangle \quad (18)$$

The screening length $\lambda_D \delta^{ab} = \Pi_{00}^{ab}(0)^{-1/2}$ leads to dissociation of bound states of a heavy quark pair, such as $(c\bar{c})$.

The suppression of J/ψ production, originally proposed by Matsui and Satz (121), is based on a simple, yet elegant idea: The ground state of $(c\bar{c})$ pair does not exist when the color screening length $\lambda_D = 1/gT$ is less than the bound state radius $\langle r_{J/\psi}^2 \rangle^{1/2}$. Lattice simulations of SU(3) gauge theory (122,123) shows that this condition should be satisfied slightly above the deconfinement temperature ($T/T_c > 1.2$). The screening length appears to be even shorter, when dynamical fermions are included in the lattice simulations (124). In addition, the D-meson is expected to dissociate in the deconfined phase, lowering the energy threshold ΔE^* for thermal break-up of the J/ψ . Blaschke(125) has estimated, using the kinetic relation

$$\sigma_{diss}^{J/\psi} \approx \sigma_0 e^{-\Delta E^*/T}, \quad (19)$$

that the dissociation probability jumps significantly already at T_c , and reaches unity at $T/T_c \approx 1.2$.

On the other hand, the J/ψ may also be destroyed in a hadronic scenario without phase transition by sufficiently energetic collisions with comoving hadrons(126), leading to dissociation into a pair of D-mesons. This mechanism has recently been analyzed carefully by Gavin(127) and by Vogt et al (128). In addition, the dependence of the suppression factor S on transverse momentum of the J/ψ is explained by a broadening of the transverse momentum distribution of projectile gluons due to rescattering(129).

6.4 Electromagnetic response function

Electromagnetic signals for the quark-gluon plasma are in many respects ideal because they probe the earliest and hottest phase of the evolution of the fireball, and are not affected by final state interactions. Electromagnetic signals probe the structure of the electromagnetic current response function(13):

$$\Pi_{\mu\nu}(q^2) = \int d^4x d^4y e^{iq(x-y)} \langle j_\mu(x) j_\nu(y) \rangle \quad (20)$$

In the hadronic phase, $\Pi_{\mu\nu}(q^2)$ is dominated by ρ^0 resonance at 770 MeV, whereas perturbative QCD predicts a broad continuous spectrum above twice the thermal quark mass $m_q = gT/\sqrt{6}$. At low $q^2 \ll 100$ MeV collective modes are predicted to exist in both phases.

The production of lepton pairs with large invariant mass in the quark-gluon plasma phase may be sensitive to pre-equilibrium phenomena, e.g. collective plasma oscillations of large amplitude(104). Such oscillations are known to occur in the framework of the chromo-hydrodynamic model where the collision energy is first stored in a coherent color field which later breaks up into $q\bar{q}$ pairs. The ensuing collective flow could enhance the production of lepton pairs of high invariant mass.

7. THE PRESENT STATUS AND FUTURE PROSPECTS

A tremendous amount of data have been gathered and analyzed, most of them from experiments with the "light" heavy ions O^{16} , S^{32} and Si^{28} . Theoretical models have been developed and widely applied to understand the implications of these data. From the available data it is easy to conclude that the temperature of about 200 MeV and energy density of about 2 GeV/fm³ are already known to have been reached in S^{32} collisions at the incident energy of 200 A GeV. Thus at the moment it can only be conceived that we do seem to produce very dense matter. But to decide whether it is 'still' hadronic matter or 'already' a quark-gluon plasma, we need more experimental information.

The exploratory phase in the still young field of QGP can be considered essentially completed. Falling short of striking discoveries, it has nevertheless provided a 'principle proof of feasibility' and even substantiated the expectation that with the next generation of experimentation we should reach new and uncharted territory. The second generation of experiments at Brookhaven, now with Au beams, is in full swing and CERN has just begun the Pb-beam program. The first aim of the new experiments should be to survey the general properties of ultra-relativistic heavy ion collisions by model calculations. Any deviation would signal the onset of a new physics. The second aim should be to explore specific probes of QGP once the general features of the reaction dynamics are understood. There is a good reason therefore to ask the question: what have we learnt from the past experiments and what we should do in the future experimental investigations to achieve these objectives?

7.1 Present Status

We discuss here what one can learn about the QGP phase transition on the basis of results available from the first generation of experiments.

7.1.1 Stopping and thermalization

First and foremost, we have learnt that heavy ions at the energies available at the AGS and SPS, about 2.5 A GeV and 10 A GeV in the c.m. system, respectively are not transparent(130). They do not appear to come to a full stop either, but this we do not know with certainty. What we can say for sure is that a considerable amount of non thermal motion (flow) is observed in the longitudinal direction. This is probably a remnant of the incident beam motion, but it is difficult to disprove that the nuclei may have come to full stop and the compressed fireball expanded longitudinally as in the Landau model. If we can trust the microscopic models, the Landau picture is excluded, though. Simple fireball models (131,132) with a longitudinal flow component have proved to be remarkably successful in describing particle yields and spectra at AGS and SPS energies. An analysis of the RQMD model (133) reveals that local equilibrium, as measured by the ratio of the spread in the transverse and longitudinal momentum distributions of hadrons $\Delta p_T / \Delta p_L$, is only achieved toward the end of the collision. At the moment of highest compression this ratio is only about one-half. At AGS energy the highest

baryon densities of roughly $10 \rho_0$ are attained (133) around 5 fm/c after first impact in Au+Au collisions. The energy density exceeds 2 GeV/fm^3 at that moment(133). These numbers agree with those obtained by applying Bjorken's formula (134)

$$\varepsilon = (\pi R^2 \tau_0)^{-1} (dE_T^{\text{max}} / dy)$$

to the (pseudo-)rapidity distributions of transverse energy(135,136) measured at CERN, which yield energy densities $\varepsilon = 2\text{-}3 \text{ GeV/fm}^3$ for S + Au collisions if $\tau_0 = 1 \text{ fm/c}$ is used for the thermalization time. Of course, there is no compelling reason for this choice of τ_0 , which at present must be considered as an arbitrary and unknown parameter.

We thus have good reason to describe the final state of heavy ion collision at the currently available energies in terms of equilibrium thermodynamic concepts, but we have equally good reason to be wary about applying these concepts to the complete duration of the collision. In particular, it appears that the hadronic matter is not quite equilibrated at the moment of highest density.

7.1.2 J/ψ suppression and strangeness enhancement

Experiment NA38 at CERN has provided us with beautiful data on J/ψ and ψ' production in nucleus-nucleus collisions (137). The data show clearly that J/ψ production is suppressed in comparison with the muon-pair continuum, and that the suppression of the ψ' is stronger than that of the J/ψ. This is, of course, exactly what is expected from a quark-gluon plasma (138). Unfortunately, it is also what is expected if the suppression is due to hadronic final-state interactions (139). The crucial question hence is whether the inelastic cross section of J/ψ with thermal hadrons is of the order of 3 mb as assumed in the hadronic models. A recent calculation (140) of the J/ψ-π cross section based on QCD sum rules claims that the cross section must be much smaller. This would rule out the hadronic mechanism and provide a strong argument in favor of a deconfined quark-gluon plasma at CERN-SPS energy. More theoretical studies and more data, especially from p + A (or A + p) collisions are needed. At present, even assertions (141) that the NA38 data are compatible with data of J/ψ suppression in p + A collisions have not been conclusively refuted.

The predicted (142) enhancement of the production of strange hadrons has been unambiguously observed at the AGS (143) as well as the SPS (144-146). Whereas the K and Λ enhancement is quite readily explained by the hadronic cascade models (147), the same models have great difficulty in describing the measured strange antibaryon yields. In fact, they fail quite miserably and have had to invoke completely new previously unobserved mechanisms, such as color ropes (148) and multi-quark clusters (149), to describe the data. On the other hand, the thermodynamical (fireball) models work extremely well (150) with or without assumption of a quark-gluon plasma, but assuming hadrochemical equilibrium. Since the dynamical models (58,52,56) at present do not provide an explanation for that assumption, we conclude that strangeness enhancement must be considered an anomalous phenomenon describing the new physics.

7.1.3 Geometry and flow

The determination of break-up geometries by intensity interferometry (HBT effect) has reached an astonishing level of sophistication (151,152). Source sizes deduced from $\pi\pi$ and KK correlations clearly exceed the radii of the projectile nuclei and agree rather well with the expected freeze out radius. $R_{\pi\pi}$ is consistently larger than R_{KK} , in agreement with expectations from the different cross sections of pions and kaons with other hadrons. The data have recently been extensively compared with predictions of the RQMD model (153), exhibiting a remarkable degree of coincidence. This gives additional confidence in these dynamical models, at least for the final stage of the collision.

The recent observation (154) of directed transverse flow in peripheral Au + Au collisions at the AGS is the first conclusive evidence for collective motion of hadrons in the final state at energies higher than those achieved at the Bevalac. Again the measurements are in reasonable agreement with the predictions of hadronic cascade models (155,156). Since the 'bounce-off' effect at high energies provides a measure of the thermal pressure in the reaction zone between the two nuclei (157), more systematic measurements may provide important information on the thermal history of the heavy ion reaction. At present, neither the HBT radii nor the observed flow are indicative of any anomalies.

7.1.4 Hadrochemistry revisited

As noted before, thermal models for the late stage of the collision have enjoyed a fair amount of success. A universal value for the freeze out temperature at AGS and SPS energies appears to be 120- 160 MeV. Deviations for different particle species can possibly be explained by the presence of transverse flow (158). Independent confirmation comes from observation of the Δ^{++} resonance in the $p\pi^+$ invariant mass spectrum in Si + Pb collisions (158,159). The Δ^{++} yield in central events is not only in remarkable agreement with RQMD predictions (157), but also agrees with freeze out temperature of 140 MeV. More evidence comes from the beautiful data of the ϕ -meson resonance in the K^+K^- channel (160).

An overall fit (138,150,151,161) yielded values for T_f , μ_B and β_T at AGS and SPS energies which have been given in the $T-\mu$ phase diagram of nuclear matter (Fig.14), from which it becomes obvious that the values of T_f and μ_B with transverse flow are very close to the putative phase boundary between hadronic gas and quark-gluon plasma. The points in the Figure correspond to the final stage of the evolution of the fireball; much more extreme conditions must be reached at earlier stages. Relying on an entropy/baryon-based argument, the authors of ref. 161 estimate initial conditions right after thermalization, which are indicated by open 'data' points in the Fig. 14. From this we cannot yet conclude that a quark-gluon plasma was formed temporarily in collisions at AGS or SPS energies, but the results are very encouraging.

7.1.5 Photons and lepton pairs

Electromagnetic probes provide information about the interior of the fireball. Experiments at CERN have looked for photons as well as for lepton pairs. Two dilepton experiments (137,162) have recently reported an excess of lepton pairs in the invariant mass region 1- 2 GeV. The WA80 experiment (163) has found evidence for direct photons from S + Au collisions in the transverse momentum range 0.5- 2 GeV/c. According to two different model calculations (164,165) these data seem to require the formation of a quark-gluon plasma which subsequently hadronizes rather slowly, forming a long-lived mixed phase with temperature 160 MeV. Unfortunately, these results are not quite conclusive, because they assume

early thermalization. It would be very interesting to see the predictions of the microscopic models (58,52) for these electromagnetic signals.

7.2 Future Prospects

It is important now to make some suggestions with the expectations in mind from the future heavy ion experiments.

7.2.1 AGS Au and CERN Pb beam program

Results from the Pb beam experiments have been reported very recently (169) by NA35, WA97, NA49 and EMU01 collaborations, but still we have to go a long way before it becomes possible to conclude something strikingly important. Obviously the most important challenge is to analyze the Au + Au data from the AGS and the new Pb + Pb data from the SPS in all detail. The large systems are much more amenable to a thermal description and also are expected to show collective behaviour more clearly. It is especially important to identify and measure the transverse flow in peripheral as well as central collisions. Strange antibaryons, charmonium, and lepton pairs presently appear to hold the most promise for anomalous behaviour. Searches for medium modifications in particle masses and widths also are important, e.g. in the vector mesons.

7.2.2 LHC and RHIC programs

(a) *Experiments*: Based on our current understanding of the data the upcoming experiments with really heavy ions at BNL and CERN should lead to the baryon densities very close to the maximum possible in any laboratory experiment, reaching or even exceeding the ones in the centre of a neutron star. The larger volumes, the (slightly) higher energy density and, most important, the increased life-time of the reaction zone will all and independently help in driving the system further towards equilibrium whether or not its internal degrees of freedom are of hadronic or partonic nature. A final quantitative interpretation of the data, and in particular a search for surviving signals from the QGP, will depend to a larger extent on a better understanding of the hadronization scenario and on the time-scales and relaxation times available in the various phases. But one should not try to hide the difficulty of the enterprise. Heavy ion collisions are complicated dynamical systems, and besides, we do not understand well enough the characteristic properties of the quark-gluon plasma to have at our disposal a series

of non ambiguous 'signatures' of its formation. Crucial to the whole program is the ability to detect the occurrence of a QGP, once it has been formed. This is not a trivial task, because the final state of ultra-relativistic heavy ion collisions is very complex having upto 10^6 produced particles, as estimated for collisions at LHC. There are two strategies for consideration:

- (i) The easiest approach is to look for observables which show such a characteristic behaviour that only a QGP may explain it ('signals'). Ideally the signal ought to be so unique that details of the dynamics should not matter.
- (ii) The possibility is to try to understand the dynamics of ultra-relativistic heavy ion collisions with aiming at a quantitative description of the reactions.

It seems at present that a pure "signal approach" (ignoring the dynamics) will not be successful. It is therefore crucial to acquire a detailed knowledge about the reaction dynamics.

Many of the proposed quark-gluon plasma signals have actually been observed, but none of these results has been demonstrated to be an unambiguous signal of the quark-gluon plasma, so far. However, one should bear in mind that the experiments were all performed with systems that were too small (Si and S are hardly "heavy" nuclei) and at energies too low to expect the formation of a full-fledged, sufficiently long-lived quark-gluon plasma state. In view of this, the experimental results presently available are encouraging.

It would be nice if the formation of quark-gluon plasma would be associated with the appearance of completely novel phenomena: there would be no ambiguity in such signatures. Indeed one should remember that the proposal to look for quark-gluon plasma in nuclear collisions (170) was originally derived from the apparent existence of unexplained phenomena observed in cosmic ray interaction, such as the famous "Centauro" events. The most probable exotic objects that might be formed from quark-gluon plasma are "stranglets" (171-173). This name describes metastable objects with baryon number $A \geq 2$ that contain several strange quarks. The simplest such object is the strangeness $S = -2$ dibaryon, the H-particle, which is predicted to be metastable in the original MIT bag model (174) and might be produced in relativistic nuclear collisions (175). Experiments (176) searching for stranglets produced in relativistic heavy ion collisions are in progress and should be encouraged. Fig. 15 gives an idea (20) about the stranglet production region.

Recently there has been speculations about the possible formation of locally "disoriented" chiral vacua in relativistic nuclear collisions (177). Such states

would decay into a large number of pions, possibly with strong isospin imbalance as observed in the Centauro events. They might be produced by some not well understood collective emission processes, or by spontaneous symmetry breaking when the dense hadronic system returns from the chirally restored phase.

(b) *Theory:* Number of theoretical presentations have been made at Quark matter '95 (178). Moreover, some model calculations have been reported in the latest literature giving interpretations of results from the completed experiments, some proposing new approaches to deal with the future investigations. Scaling behaviour in the string fusion model for heavy ion collisions (179), rapid hadronization scenario based on the dynamical chiral model (180), formation time of hadrons and dense matter produced in relativistic heavy ion collisions (181), nucleation of QGP (182), mass dependence of hadronization (183), and HBT interferometry of expanding source (184) have been discussed. So far as other signals are concerned the hard processes in the hadronic interactions have been exclusively studied (185). Moreover, calculations for particle production and strangeness enhancement (186-190) as well as diphoton production (191) have been discussed. However, there is some confusion regarding models explaining the experimental data. For example, thermal as well as non thermal models (e.g. RQMD, DPM string fusion) all explain data reasonably well and one cannot decide which approach should be considered as more appropriate for QGP investigations. Thus more detailed analysis of the dynamical contents of state-of-the-art hadronic cascade models is urgently needed. Are they self-consistent, or do they contain internal inconsistencies, such as regions with overlapping hadrons indicative (165) of quark matter formation? Such studies can also help to answer the question how good thermal models really are and at what stage of the collision they can be applied. Models of the freeze-out phase should begin to consistently include variations in the mean free paths of different hadrons (166). Finally, there is a real need for more sophisticated dynamical models of the transition from a quark-gluon plasma to a hadronic gas. (167-168).

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Figure 1. Thermal history of the Universe

Figure 2. Two-chain diagram describing particle production in high energy nucleon-nucleon collision due to (a) color exchange in DPM and (b) momentum exchange in Lund Model Fritiof.

Figure 3. The phase diagram of strongly interacting matter, showing the hadronic phase at low temperature and baryon density, the transition region (mixed phase) and the QGP phase. The full curves lines illustrate trajectories followed in supernovae explosions, Big Bang evolution, and possibly in heavy-ion reactions at present and future accelerators.

Figure 4. Schematic representation of a heavy-ion collision at impact parameter b , assuming 'clean cut' geometry.

Figure 5. A collision of two large nuclei $A + B$ in a space-time diagram. Here the two nuclei collide along the light-cone trajectories marked A and B. Particles are produced in the region of the forward light cone. The various stages of the expanding matter at various times are indicated in the figure.

Figure 6. Stability line of the quark-gluon plasma in the $T-\mu$ plan for two values of the strong coupling constant. The lines indicate where the pressure of a gas of quarks and gluons vanishes. The quark-gluon plasma is unstable in the lower left region of the figure.

Figure 7. Energy density and pressure of a pion gas and a free quark-gluon gas as function of temperature. A first-order phase transition occurs where the pressure of the two phases is equal. The solid line indicates the stable phase.

Figure 8. Energy density ϵ and pressure P of pure glue matter as calculated by simulations of SU(3) lattice gauge theory (96). ϵ and P are plotted relative to the Stefan-Boltzmann limit, Eq (6a)

Figure 9. QCD diagrams describing scattering processes contributing to thermalization of the quark-gluon plasma.

Figure 10. String-based picture of the formation of a quark-gluon plasma. Primary interactions lead to color flux tubes which break by quark pair production.

Figure 11. Schematic view of a parton cascade

Figure 12. Rapidity distributions of charmed quarks produced in secondary gluon interactions, in comparison to the primary parton model predictions.

Figure 13. Average transverse momentum of emitted hadrons as function of transverse energy dE_T/dy , representing the maximum energy density reached in a collision. The different curves correspond to: (a) pion gas, (b) Hagedorn resonance gas, (c) quark-gluon plasma.

Figure 14. Thermal freeze-out parameters shown in the phase diagram of nuclear matter. Two lines are shown to indicate the degree of uncertainty. Horizontal axis shows the quark chemical potential $\mu_q = 1/3 \mu_B$.

Figure 15. Phase diagram of nuclear and quark matter plotted as functions of the logarithm of the chemical potential. As shown there are regions of the plot where matter resides in the strange quark phase as a dilute gas of stranglets. In regions of very high baryochemical potential, strange quark matter would be found in bulk.