Matter & More in Nuclear Collisions

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Abstract:

The aim of high energy nuclear collisions is to study the transition from hadronic matter to a plasma of deconfined quarks and gluons. I review the basic questions of this search and summarize recent theoretical developments in the field.

1. New States of Matter

Statistical QCD predicts that high temperatures and baryon densities will lead to new states of strongly interacting matter. Increasing $T$ at low baryon density transforms a meson gas into a deconfined plasma of quarks and gluons (QGP); this transition has been studied extensively in computer simulations of finite temperature lattice QCD. High baryon densities at low $T$ are expected to produce a condensate of colored diquarks. The resulting phase diagram in terms of temperature and baryochemical potential $\mu$ is schematically illustrated in Fig. 1, with hadronic matter as color insulator, the QGP as color conductor, and the diquark condensate as color superconductor.

![Phase diagram of strongly interacting matter.](image)

Figure 1: Phase diagram of strongly interacting matter.

With high energy nuclear collisions, we want to study in the laboratory the deconfinement transition and the properties of the QGP. Hard probes, such as the production of quarkonia, open charm and beauty, jets and photons are expected to provide information

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about the hot early stages of the produced medium. The masses and decays of different hadrons, their momenta, correlations and relative abundances constitute the soft probes to study the later stages of the medium and its freeze-out.

In this concluding talk, I will summarize some recent theoretical developments in the field, without any claim to completeness. My emphasis will be on concepts more than on specific models, and on questions more than on answers.

2. Thermalization

Since the basic purpose of the experimental program is to produce strongly interacting matter, it is of central importance to determine if, how and when the non-thermal initial state of two colliding nuclei becomes thermalized.

In a nuclear collision, the incoming nucleons or their secondaries can interact with the other target nucleons. This results in nuclear phenomena such as the Cronin effect, normal nuclear quarkonium suppression or parton energy loss in normal nuclear matter. These effects do not involve any new produced medium. To achieve that, the secondaries coming from different sources must interact, a phenomenon referred to as color interconnection, exogamous behavior or cross talk [1].

A test for such cross-talk has been considered in $e^+e^-$ annihilation into hadrons at $\sqrt{s} = 2M_W$, as studied in LEP experiments at CERN [2]. The reaction first leads to $W^+W^-$ production; subsequently, one possibility is that each of the two $W$'s decays into a $q\bar{q}$ pair, which then hadronizes. An alternative channel has one of the two $W$'s undergo leptonic decay into a neutrino and a lepton. In the case of cross talk, it is predicted [1] that for the resulting hadron multiplicities, $N_h(q_1\bar{q}_1, q_2\bar{q}_2) \neq 2N_h(q\bar{q}, \nu l)$; in addition, the source radii obtained through HBT studies should be different in the two channels. Neither of these predictions is supported by LEP data, so that so far there is no evidence of cross talk shown by the hadrons produced in high energy $e^+e^-$ annihilation. This also excludes cross talk at earlier partonic stages.

Do $AA$ collisions with their much higher density of superimposed interactions lead to cross talk? We address this question by looking at hadron abundances. It is found that these are quite well described by the predictions of an ideal resonance gas [3], parametrized in grand-canonical form by a freeze-out temperature $T_f$ and a baryochemical potential $\mu_B$. With increasing collision energy $T_f$ converges to about 170 MeV (Fig. 2). This alone does not, however, allow us to conclude that we have indeed obtained a thermal system with full cross-talk. It is known [3] that also the elementary hadroproduction processes initiated by $e^+e^-$ annihilation or $pp/p\bar{p}$ scattering lead to thermal hadron abundances, with freeze-out temperatures which agree very well with those observed in $AA$ collisions; they are included in Fig. 2. There is, however, one important difference: in the elementary reactions, there are fewer strange hadrons than predicted by a grand-canonical resonance gas, while $AA$ collisions do not show such a strangeness reduction. The change in relative strangeness production, referred to as strangeness suppression or enhancement, depending on the point of view, has found a very natural explanation in the observation [3] that in elementary processes, with rarely more than one $s\bar{s}$ pair per interaction region, strangeness conservation has to be taken into account exactly and not just “on the average”, as implied in a grand-canonical formulation. It had in fact been observed long ago that the exact local conservation of quantum numbers can reduce the relative production rates by several orders of magnitude in comparison to the average grand-canonical rates [3].
The transition from exact to average (grand-canonical) strangeness conservation does imply, however, that the $AA$ collisions behave as one large system, not as a sum of many elementary collisions. The strange hadrons produced in one elementary collision of an $AA$ interaction must be aware of the strange hadrons produced in other $NN$ collisions, if a grand-canonical description is valid. So at least at the hadronization stage of the medium produced in nuclear collisions, there is cross talk – it makes sense to speak of a large-scale hadronic medium.

One remaining question in this context is the correct treatment of hidden strangeness; if the $s\bar{s}$ nature of the $\phi$ is not taken into account, its production rates are quite generally overpredicted.

In addition, there remains the tantalizing question of why the hadron abundances in elementary processes like $e^+e^-$ or $pp/p\bar{p}$ already follow the pattern predicted by an ideal resonance gas. It might just indicate that the non-perturbative hadronization of quarks and gluons simply proceeds such as to maximize the entropy: partons hadronize into the states of a resonance gas with largest phase space. The associated temperature is then the limiting temperature of such a system, the Hagedorn temperature, and hence is universal.

Since the initial structure of e.g. $e^+e^- \rightarrow$ hadrons is that of a fast $q\bar{q}$ pair emitting gluons which later hadronize, the annihilation process is certainly not thermal in its early stages. Thermal hadron abundances thus do not imply that the previous partonic system was already thermal.

Partonic cross talk is the basic input of most parton cascade models and thus has been considered in the corresponding codes for quite some time. Its role in the establishment of parton thermalization has recently been addressed in an interesting conceptual treatment [8]. The two incoming nuclei can be viewed in the central rapidity region (where the valence quarks are unimportant) as gluon beams in which each gluon has a transverse radius $r_g \sim 1/k_T$ determined by its transverse momentum $k_T$. The geometric interaction cross section for gluon of sufficient hardness, $k_T \geq q_0$, is $\sigma_g \sim \alpha_s(q_0)r_g^2$. We would like to know when these gluons of sufficient hardness begin to overlap in the transverse area determined by the nuclear radius $R_A$. Full cross talk evidently occurs when a gluon on one side of the nuclear disk $\pi R_A^2$ is connected by overlapping gluons to a partner on the
other side of the disk. This starts at the percolation point

\[ N_g \sigma_g \simeq \pi R_A^2, \]  

(1)

where \( N_g(k_T) = A x g(x) \) denotes the number of gluons as determined from the gluon distribution function \( g(x) \) at central Bjorken \( x \sim \sqrt{s} \). From Eq. (1) we thus obtain

\[ q_s^2 \sim \alpha_s A^{1/3} x g(x) \]  

(2)

for the saturation momentum \( q_s \). When \( k_T \gg q_s \), the gluons form a dilute and hence disjoint system in the transverse plane, without cross talk. At \( k_T = q_s \) percolation sets in and the gluons form a connected interacting system, with full cross talk. For such a system, one can estimate that thermalization occurs after a time

\[ \tau_0 \sim \alpha_s^{-13/5} q_s^{-1}, \]  

(3)

leading to a thermal gluon medium of temperature

\[ T_0 \sim \alpha_s^{2/5} q_s. \]  

(4)

From deep inelastic scattering one has \( x g(x) \sim (\sqrt{s})^\lambda \), with \( \lambda \simeq 0.2 \). Together with Eq. (2), this implies that large nuclei (large \( A \)) and/or high collision energies (large \( \sqrt{s} \)) lead to early parton thermalization and a hot QGP. A very important task for theory is clearly to turn these conceptual considerations into a quantitative formalism.

### 3. Hadrons in Matter

Do \( AA \) collisions produce an interacting hadronic medium, or does the earlier partonic state hadronize directly into an ideal resonance gas? That is the main question to be addressed here. It is of particular interest in view of chiral symmetry restoration. For temperatures \( T < T_c \), the massless quarks of the QCD Lagrangian \( \mathcal{L}_{\text{QCD}} \) “dress” themselves through gluon interactions to become constituent quarks with an effective mass \( M_q \sim 0.3 \) - 0.4 GeV, thereby spontaneously breaking the chiral symmetry of \( \mathcal{L}_{\text{QCD}} \). For \( T = T_c \), chiral symmetry is restored and \( M_q \to 0 \); we have here assumed a system of vanishing baryon number density, for which deconfinement and chiral symmetry restoration coincide at \( T = T_c \). Since the vector meson mass \( M_\rho \simeq 2 M_q \), the behavior of \( M_\rho(T) \) in an interacting hadronic medium for \( T \to T_c \) would be a way to study the onset of chiral symmetry restoration.

The in-medium behavior of hadron masses can be calculated in finite temperature lattice QCD. First studies addressed the temperature dependence of the screening mass in quenched QCD. Below \( T_c \), they showed very little \( T \)-dependence; but in view of the noted simplifications, they are presumably not really conclusive. Today it is possible to calculate the actual pole mass, but with present computer performance still only for the quenched case \([9]\). The advent of more powerful computers in the next 2 - 3 years should, however, lead to such calculations in full QCD with light quarks. The present quenched studies show that above \( T_c \), there are no more mesonic bound states; scalar and pseudoscalar correlations agree, indicating chiral symmetry restoration. Below \( T_c \), the results for the pole masses also show rather little \( T \)-dependence; however, this may well be an artifact of quenching, as the following considerations seem to indicate.
Lattice studies of the heavy quark potential $V_Q(T, r)$ in full QCD (with light dynamical quarks) show at all temperatures $T$ string breaking in the large distance limit $r \to \infty$. At $T = 0$, the string connecting two heavy color charges should break when its energy surpasses the mass of a typical light hadron, or equivalently when

$$V_Q(T, r = \infty) \simeq 2M_q(T).$$

Hence $M_q(T)$ can be determined in finite temperature lattice calculations of $V_Q(T, r)$. Note, however, that the constituent quark mass here is obtained from a heavy-light meson and could contain some dressing of the heavy quark; hence it need not coincide fully with that from a light-light meson such as the $\rho$. Heavy quark potential studies have recently been carried out for $N_f = 2$ and 3 for a range of different quark masses $[10]$. They show in particular that

- for $m_q \lesssim 0.4 T \simeq 60$ MeV, the dependence of $V_Q$ on $m_q$ becomes negligible, indicating that the chiral limit is reached;
- string breaking occurs earlier (at smaller $r$) with increasing temperature (Fig. 3);
- from $V_Q(T = 0, r = \infty) \simeq 1$ GeV it follows that $M_q(T = 0) \simeq 0.5$ GeV, while for $T \to T_c$, $V_Q(T)$ and hence also $M_q(T)$ vanish.

The resulting temperature dependence of the heavy quark potential is shown in Fig. 3. It is seen that the approach of chiral symmetry restoration leads to a pronounced variation of $V_Q(T, \infty)$ and hence of $M_q(T)$ with $T$. This suggests that also the mass of the $\rho$-meson should show such a temperature variation, in contrast to the present pole-mass results from quenched lattice QCD. A study of the temperature dependence of hadron masses in unquenched lattice calculations would thus be of great interest.

4. Partons in Matter

Since the energy loss of a fast parton passing through a medium will depend on the nature of this medium, jet quenching should provide a tool to specify the state of matter
produced in nuclear collisions. At lower momenta, the energy loss can occur through ionisation of the constituents of the medium; at high momenta, gluon radiation of the passing parton is the main mechanism.

The crucial feature in radiative energy loss is the formation time $t(k)$ or the formation length $z(k)$ for a gluon of momentum $k$, compared to the intrinsic scales of the medium: the mean free path $\lambda$, the mean distance $d$ between scattering centers and the overall linear size $L$ of the medium. For $\lambda > d > z(k)$, the radiated gluons see independent charges and the scattering is incoherent. For $d < \lambda < z(k)$, there is coherent scattering of the nascent gluon with several scatterers, leading to destructive interference; this is the so-called Landau-Pomeranchuk-Migdal (LPM) effect which reduces the energy loss. In Fig. 4 we compare schematically the incoherent form $dE/dz \sim -E$, where $E$ is the parton energy, to the LPM form $dE/dz \sim -\sqrt{E}$. In a medium of small linear size, $L < L_c$, there is a further finite size reduction, leading to

$$- \frac{dE}{dz} \simeq \frac{3\alpha_s}{\pi} \left\{ \begin{array}{ll}
(\mu^2 E/\lambda)^{1/2}, & L > L_c \\
(\mu^2/\lambda) L, & L < L_c
\end{array} \right.$$  \(6\)

for the energy loss in a quark-gluon plasma. Here $L_c \equiv (E\lambda/\mu^2)^{1/2}$ denotes the limiting length scale and $\mu^{-1}$ the screening length of the medium.

Figure 4: Parton energy loss through coherent vs. incoherent scattering (left) and through coherent scattering in a medium of finite size (right).

For a parton traversing a QGP at temperature $T = 250$ MeV over a length of $L = 10$ fm, Eq. (6) leads to an energy loss of about 30 GeV; this is to be compared to indications that cold nuclear matter of the same size would only result in an energy loss of 2 GeV \[12\]. In contrast to the QGP results, the value for a normal nuclear medium is really only an estimate, and thus only jet production data from $pA$ collisions can provide a reliable basis for comparison.

Another problem is immediately evident from the QGP calculation. What transverse momenta are really needed to specify a jet? Present RHIC data stop around at leading particles of some 5 GeV, which presumably is well below the value for the jets assumed in the QCD studies.

We add a comment here on the interpretation of the RHIC results on high $p_T$ hadrons. The measured spectra have to be compared to some reference spectrum in order to look for a possible quenching, and this is generally based on binary collisions: the spectra
from central AA collisions are normalized to pp (or peripheral AA) data multiplied by the number of binary collisions. The resulting ratio (see Fig. 5) is well below unity in the entire range $0 \lesssim p_T \lesssim 5$ GeV. If one would instead use the number of wounded nucleons as reference, the corresponding ratio will be larger than unity for almost all values of $p_T$. From multiplicity studies it is clear that the overall data, dominated by low to intermediate $p_T$, are well below what is expected from binary collision scaling. Hence in order to obtain a reliable reference, one should use an interpolating form of the type

$$\left( \frac{dN}{dp_T^2} \right)_{AA}^\text{ref} = \left[ N_w \left( 1 - \frac{p_T^2}{a + p_T^2} \right) + N_c \left( \frac{p_T^2}{a + p_T^2} \right) \right] \left( \frac{dN}{dp_T^2} \right)_{pp},$$

(7)

where $N_w$ denotes the number of wounded nucleons and $N_c$ the number of binary collisions; the parameter $a$ determines the relative importance of the two types of production mechanisms. Instead of being set to zero, as in present studies, it should be choses such as to correctly reproduce the measured multiplicity. This is expected to lead to a behavior like that shown in Fig. 5, with a Cronin-like pattern at relatively low $p_T$ followed by a suppression below unity, and one could then clearly define quenching effects.

The last point to be addressed in this section concerns the radiative energy loss of heavy quarks traversing a QGP. It was noted [13] that for massive quarks the gluon emission suffers a ‘dead-cone’ effect, which suppresses radiation for forward angles

$$\theta \lesssim M_Q/\sqrt{P_Q^2 + M_Q^2},$$

(8)

where $M_Q$ denotes the mass of the heavy quarks. This radiation suppression in turn reduces the energy loss of heavy quarks and thus predicts an increase of the ratio $D/\pi$ for high $p_T$.

![Figure 5: Expected transverse momentum distribution from AA collisions normalized to a reference distribution interpolating from a wounded nucleon to a binary collision model (see Eq. 7).](image)

5. Quarkonia in Matter

The essential feature that distinguishes the quarkonium ground states $J/\psi$ and $\Upsilon$ from the normal light hadrons is their much smaller radius (about 0.2 fm for the $J/\psi$ and about
0.1 fm for the \( \Upsilon \), due to the much higher bare quark mass \( m_c \approx 1.4 \text{ GeV}, m_b \approx 4.5 \text{ GeV} \). Their binding is thus largely due to the Coulombic part of the QCD potential \( \sigma r - \alpha/r \); the string tension \( \sigma \) does not matter very much. Equivalently, the gluon dressing which makes massive constituent quarks out of the almost massless light quarks (with \( M_q \approx 0.3 - 0.4 \text{ GeV} \)) has little effect on the heavy quarks. In a medium approaching the deconfinement point, for \( T \to T_c \), the string tension vanishes: \( \sigma(T) \to 0 \), as does the constituent quark mass because of chiral symmetry restoration, \( M_q \to 0 \). As a consequence, light and light-heavy hadrons disappear, but sufficiently tightly bound quarkonia will persist even above \( T_c \) and can thus serve as probes of the quark-gluon plasma. These arguments also suggest that quarkonium masses decrease less with temperature than the masses of the open charm or beauty mesons \( D \) and \( B \). As a consequence, the open charm threshold can in a hot medium fall below the mass of previously stable higher excited charmonium states and thus allow their strong decay, and similarly for bottomonia.

We therefore want to compare \( 2M_D(T) \) with \( M_i(T) \), where \( i \) specifies \( \psi', \chi \) and \( J/\psi \), as well as the corresponding \( b \)-quark states. For the masses of the open charm/beauty states, we make use of the lattice studies already introduced in section 3. The string breaking potential introduced there determines with \( V(T, r = \infty) \approx 2M_q(T) \approx 2(M_D - m_c) \) effectively the \( D \)-mass. The quarkonium masses can be obtained by solving the Schrödinger equation with the potential \( V(r, T) \) determined in the same lattice studies. Comparing the temperature dependence of the light-heavy masses to that of the quarkonium states shows two distinct types of behavior [14, 15].

In Fig. 6 we see that with increasing temperature \( 2M_D \) and \( 2M_B \) indeed drop below the masses of the highest excited states, \( \psi' \) and \( \chi_c \) for charmonia, \( \Upsilon^* \) and \( \chi_b' \) for bottomonia, respectively, before the deconfinement point is reached. These states thus disappear in a hot hadronic medium through in-medium decay into open charm/beauty. If the dissociation thresholds are experimentally measured, they thus specify the temperature of the hot but still confined system at four different points, tracing out the approach of chiral symmetry restoration.

The mass gaps of \( J/\psi, \Upsilon, \chi_b, \chi_c \) and \( \Upsilon \) at \( T = 0 \) are much larger than \( \Lambda_{\text{QCD}} \), so that we expect them to survive deconfinement and be dissociated only by color screening in the quark-gluon plasma, as originally proposed for the \( J/\psi \) [16]. This dissociation sets in when the intrinsic scale of the quarkonium, its radius, falls below the screening radius as the scale characterizing the medium. In Fig. 6 this effect is seen to occur for the \( \Upsilon \) at \( T \approx 2.3 T_c \). The other mentioned states persist up to about \( T_c \) when compared to open charm/beauty masses (see Fig. 3); in a screening approach, they are dissociated in a QGP just slightly above \( T_c \) (see Fig. 7 for the \( J/\psi \)). Given the accuracy of the present lattice results near deconfinement, and in view of the possible break-down of a Schrödinger equation near \( T_c \), we can thus only conclude that \( J/\psi, \chi_b \) and \( \Upsilon \) are dissociated approximately at \( T_c \).

We thus obtain a thermal quarkonium dissociation pattern which indeed is very similar to that provided by the spectral lines from stellar matter [18]. The suppression thresholds for \( \psi', \chi_c \), \( \Upsilon^* \) and \( \chi_b' \) specify a hot hadronic medium different temperatures; when the \( J/\psi \), \( \Upsilon \) and \( \chi_b \) disappear, deconfinement is reached, and the dissociation point of the \( \Upsilon \) indicates a hot QGP, with \( T > 2 T_c \).

The observed production of \( J/\psi \) and \( \Upsilon \) occurs in part through the decay of higher excited states, such as \( \chi_c \to J/\psi \); the respective fractions of the different contributions are known experimentally or can be determined from data [15]. Since such decays take place
far outside the interaction region, the produced medium sees and suppresses the different ‘parent’ states. This leads to the well-known sequential suppression pattern \[17, 15\] which distinguishes thermal threshold behavior such as deconfinement form dissociation by hadronic comover scattering (see \[19\] for a survey). Perhaps the most interesting feature which has so far emerged from nuclear collision studies is the observation of just such a multistep structure of \(J/\psi\) suppression \[20\]. Future experiments, both at CERN (NA60) and at RHIC, will undoubtedly provide further details to check if this structure is indeed due to sequential quarkonium suppression.

Our considerations so far have ignored possible fluctuations of the medium. To illustrate, we note that for \(T > 0\) the mass of the \(D\) and to a lesser extent the mass of the \(\chi_c\) will fluctuate around the values we have here calculated. Instead of a \(\delta\)-function, we will have a peak with a certain width (collision broadening), which in principle can be provided by lattice calculations. For a conclusive study of sequential quarkonium suppression in nuclear collision this would seem a prerequisite.
A second open question concerning applications to experiment is a reliable determination of the energy densities or temperatures attained there. All present lattice studies find $T_c \simeq 0.15 - 0.20$ GeV for the deconfinement temperature; the corresponding energy density is $\epsilon(T_c) \simeq 1$ GeV/fm$^3$, although it then grows quickly to values near the Stefan-Boltzmann limit, so that $\epsilon(1.1 T_c) \simeq 2$ GeV/fm$^3$. Presently quoted values for the energy densities in $Pb - Pb$ collisions at the CERN-SPS are in the range 2 - 3.5 GeV/fm$^3$; they are based on Bjorken’s estimate, which for central collisions gives

$$\epsilon \simeq \left( \frac{dN_h}{dy} \right)_{y=0} \frac{p_0}{\pi R_A^2 \tau_0},$$

(9)

where $dN_h/dy$ denotes the multiplicity and $p_0$ the average energy of the produced hadrons, $R_A$ the nuclear radius and $\tau_0 \simeq 1$ fm some average formation time of the medium. Obviously the choice of $\tau_0$ is rather crucial, and a cross check of the reliability of the resulting estimates would thus seem very necessary.

Summary

- Hadron abundances, in nuclear collisions as well as in elementary interactions, follow the pattern of an ideal resonance gas. The strangeness suppression observed in elementary processes appears accountable through exact strangeness conservation. The observed energy independent freeze-out temperature $T_f \simeq 170$ MeV seems to reflect critical features.

- Finite $T$ lattice studies of the heavy quark potential $V(T, r)$ show a significant variation of the string breaking energy $V(T, \infty)$ for $T \rightarrow T_c$. This could be an indication for a similar temperature variation of light hadron masses in full QCD.

- Fast partons passing through a QGP suffer a considerable energy loss, which should be observable for sufficiently hard jets or their decay products. Present RHIC data require a reference distribution interpolating from a wounded nucleon to a binary collision form.

- Quarkonia in hot matter can be dissociated by two distinct mechanisms. Higher excited states decay strongly into open charm/beauty mesons when the masses of the latter decrease as the system approaches chiral symmetry restoration. More tightly bound lower states survive up to deconfinement and are subsequently dissociated by color screening in the hot QGP.

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