COLOUR DECONFINEMENT
IN HOT AND DENSE MATTER*

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Summary:

We first introduce the conceptual basis of critical behaviour in strongly interacting matter, with colour deconfinement as QCD analog of the insulator-conductor transition and chiral symmetry restoration as special case of the associated shift in the mass of the constituents. Next we summarize quark-gluon plasma formation in finite temperature lattice QCD. We consider the underlying symmetries and their spontaneous breaking/restoration in the transition, as well as the resulting changes in thermodynamic behaviour. Finally, we turn to the experimental study of strongly interacting matter by high energy nuclear collisions, using charmonium production to probe the confinement status of the produced primordial medium. Recent results from Pb-Pb collisions at CERN may provide first evidence for colour deconfinement.

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1. Critical Behaviour in Strongly Interacting Matter

At sufficiently high temperatures or densities, matter is expected to undergo a transition from confined to deconfined quarks as its basic constituents. It is, according to quantum chromodynamics (QCD), always made up of quarks and gluons; in the confined phase, these coloured constituents are bound to colour-neutral hadrons, while in the deconfined quark-gluon plasma (QGP), there are freely moving colour charges.

We begin with a look at the critical behaviour expected during the transition from hadronic matter to QGP. The Coulomb potential between two electric charges becomes Debye-screened in a medium of many other charges, reducing its range to the Debye radius $r_D$,

$$
\frac{e}{r} \rightarrow \frac{e}{r_D} e^{-r/r_D}.
$$

(1)

At high enough density, when the Debye radius becomes shorter than the atomic radius, bound electrons are liberated into the conduction band, changing an insulator into a conductor. Similarly, colour charges bound by the linearly rising confinement potential of QCD become screened in a dense medium,

$$
\sigma r \rightarrow \sigma r_c [1 - e^{-r/r_c}],
$$

(2)

with $r_c$ as colour screening radius. At sufficiently high density, colour screening will therefore dissolve a hadron into its coloured quark constituents, so that deconfinement is the QCD analog of the insulator-conductor transition.

When the electrons of an insulator are decoupled, their mass is shifted from the standard $m_e$ to an effective value $m_{\text{eff}}$, determined by lattice interactions and the effect of the electron gas in the conductor. In QCD, we expect the quark mass, which takes on an effective constituent quark value $m_{q}^{\text{const}} \approx m_{\text{proton}}/3$ when the quark is confined to a hadron, to drop back to the current quark value $m_q$ of the QCD Lagrangian once it is no longer confined. Such a quark mass shift is therefore another aspect to be considered in the course of colour deconfinement. In the limit $m_q = 0$, the Lagrangian becomes chirally symmetric, so that in this case the intrinsic chiral symmetry of the theory must be spontaneously broken in the hadronic phase and restored in the QGP.

The basic condition for the transition from hadronic to quark matter is a sufficiently high density of constituents: it then becomes impossible to define a given quark-antiquark pair or a quark triplet as some specific hadron, since within any hadronic volume there are many other possible partners. Such a density can be achieved either by compressing baryons (cold nuclear matter) or by heating a mesonic medium, increasing its density by particle production in collisions (hot mesonic matter). The phase diagram of QCD can thus maps out regions in the plane of temperature $T$ and baryochemical potential $\mu$ (see Fig. 1), with the latter specifying the mean baryon number density (baryons minus antibaryons).

How many phases are there in QCD thermodynamics? As we shall see shortly, it is known from lattice QCD studies that for $\mu = 0$, deconfinement and the approximate chiral symmetry restoration associated to light quarks occur at the same temperature $T_c$. So far, technical reasons prevent us from carrying out lattice calculations for $\mu \neq 0$,
leaving in particular a “terra incognita” in the low temperature, high density region. Since the potential between quarks contains an attractive component, it is quite conceivable that after deconfinement there will be diquark formation, with quark pairs playing the role of Cooper pairs in a diquark phase similar to a superconductor. Only at high enough density or temperature, such diquarks would then break up to form the true QGP. It appears that this really interesting question of the low temperature structure of QCD matter will have to remain unanswered until a suitable lattice scheme is developed for baryonic matter.

Fig. 1: The phase structure of strongly interacting matter

It is obviously of great interest to estimate the hadron-quark transition temperature, and this can be done in various phenomenological models – bootstrap model, bag model, string model, dual resonance model or even percolation theory. It is quite reassuring that they all lead to very similar values. Let us consider the perhaps simplest picture by assuming the hadronic phase to be an ideal gas of massless pions, the quark phase an ideal gas of massless quarks and gluons, based on colour SU(3) and the two light quark flavours. The pressure of the former is

$$P_h = 3 \frac{\pi^2}{90} T^4 \approx \frac{1}{3} T^4,$$  \hspace{1cm} (3)

taking into account the three charge degrees of freedom of a pion. For the quark-gluon system we get

$$P_q = [2 \times 8 + \frac{7}{8} (2^3 \times 3)] \frac{\pi^2}{90} T^4 - B \approx 4 T^4 - B,$$ \hspace{1cm} (4)

with two spin and eight colour degrees of freedom for the gluons, two spin, two flavour, two particle-antiparticle and three colour degrees for the quarks. The bag pressure $B$ characterizes the difference between the physical vacuum and the ground state of QCD;
it provides a model for the confining feature of the theory, forcing the quarks into a hadronic volume. Since the preferred thermodynamic state is that of highest pressure (lowest free energy), the cross-over point $T_c$ obtained by equating $P_h$ and $P_q$ defines the critical temperature

$$T_c = \left( \frac{90 B}{34 \pi^2} \right)^{1/4} \simeq 0.7 B^{1/4}$$

separating the low temperature hadron from the high temperature QGP phase, with a transition which is by construction of first order. Using a bag pressure value from charmonium spectroscopy ($B^{1/4} \simeq 0.2$ GeV), we find with $T_c \simeq 140$ MeV the Hagedorn temperature first obtained through the statistical bootstrap model.

After these phenomenological preliminaries, let us now see what results about deconfinement can be obtained directly from statistical mechanics based on QCD as underlying theory.

2. Statistical QCD

QCD as the dynamical theory of strong interactions is defined by the Lagrangian

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu}^a F_{\mu\nu}^a - \sum_f \bar{\psi}_f^i (i\gamma^\mu \partial_\mu + m_f - g \gamma^\mu A_\mu)^{\alpha\beta} \psi_f^j ,$$

with

$$F_{\mu\nu}^a = (\partial_\mu A_\nu^a - \partial_\nu A_\mu^a - g f_{bc}^a A_\mu^b A_\nu^c) .$$

Here $A_\mu^a$ denotes the gluon field of colour $a$ ($a=1,2,\ldots,8$) and $\psi_f^i$ the quark field of colour $\alpha$ ($\alpha=1,2,3$) and flavour $f$; the current quark masses are given by $m_f$. With the dynamics thus fixed, we obtain the corresponding thermodynamics from the partition function, which is most suitably expressed as a functional path integral,

$$Z(T,V) = \int dA d\psi d\bar{\psi} \exp \left( -\int_V d^3x \int_0^{1/T} d\tau \mathcal{L}(A,\psi,\bar{\psi}) \right) ,$$

since this form involves directly the Lagrangian density defining the theory. The spatial integration in the exponent of Eq. (8) is performed over the entire volume $V$ of the system; in the thermodynamic limit it becomes infinite. The time component $x_0$ is “rotated” to become pure imaginary, $\tau = ix_0$, thus turning the Minkowski manifold, on which the fields $A$ and $\psi$ are originally defined, into a Euclidean space. The integration over $\tau$ in Eq. (8) runs over a finite slice whose thickness is determined by the temperature of the system. The finite temperature behaviour of the partition function in the Euclidean form thus becomes a finite size effect in the imaginary time direction. Eq. (8) is derived from the usual trace form of the partition function. As a consequence, vector fields have to be periodic and spinor fields antiperiodic at the boundaries of the imaginary time integration.

Once $Z(T,V)$ is given, we can calculate all thermodynamical observables in the usual fashion. Thus

$$\epsilon = (T^2/V) \left( \frac{\partial \ln Z}{\partial T} \right)_V$$
gives us the energy density, and

\[ P = T \left( \frac{\partial \ln Z}{\partial V} \right)_T \]  

the pressure. What remains is to find a way to actually evaluate these expressions in the case of a relativistic, interacting quantum field theory. In QED, there are divergences both for small (infrared) and for large (ultraviolet) momenta; hence renormalization is required to get finite results, and it is to be expected that renormalisation will be necessary for QCD as well. But there is a further, more serious, problem. The standard evaluation method for QED – perturbation theory – is not applicable to the study of critical behaviour. Since long range correlations and multi-particle interactions are of crucial importance here, the interaction terms cannot be assumed as small. We therefore need a non-perturbative regularisation scheme for the solution of a relativistic quantum field theory. So far, there is only one method available which fulfills these requirements: the lattice formulation introduced by K. Wilson [1]. It puts the thermodynamic observables, such as the energy density (9) or the pressure (10), into a form that can be evaluated numerically by computer simulation [2].

The lattice formulation of statistical QCD is obtained as follows. First, the continuum integration over space and (imaginary) time in the action (the exponent in Eq. (8)) is replaced by a summation over a finite space-time grid. To maintain a gauge invariant form, the gluon fields must be associated to the links between adjacent sites of this lattice, the quark fields to lattice sites. Next, the resulting discrete form of the action allows the quark field integration in Eq. (8) to be carried out. This leads to the partition function

\[ Z(T,V) = \int \prod_{\text{links}} dU \exp[-S(U)], \]  

where the \( U \) are unitary matrices formed from the gluon fields. The temperature is determined by the number \( N_t \) of lattice sites in the imaginary time direction, \( T = 1/N_t a \), the volume by the corresponding number in space, \( V = (N_s a)^3 \), with \( a \) denoting the lattice spacing. The action \( S(U) \) in Eq. (11) is found to have the form of a (gauge-invariant) spin system, so that the partition function becomes structurally equivalent to that of a spin system. And for the study of such systems, there are known computer simulation methods.

In computer simulation one essentially creates a lattice world according to the given dynamics (here QCD) on a large computer and brings this world into equilibrium by successive Monte Carlo iterations, using the action \( \exp\{-S(U)\} \) as a weight to determine improved configurations. Once equilibrium is obtained, one measures any observable of interest on a large number of equilibrium configurations and thus obtains its value (for a more extensive survey, see [3]). In this way, one can determine the behaviour of the energy density, the specific heat, or any other desired quantity even in the region of critical behaviour of the system [4].

However, the method does encounter some problems. We are evidently interested in true statistical QCD as a continuum theory, not in its approximation on a discrete
lattice. One thus has to study how lattice results behave in the limit of large lattice size and small lattice spacing. This extrapolation to the continuum limit requires the study of lattices of different (large) sizes, so that extensive numerical work on large-scale computers is needed. Such work has been going on over the past fifteen years, but for really precise results of the full theory with quarks, more work is needed. In particular, the smaller the quark mass is, the more time-consuming the calculations become. Hence present results still use somewhat too large current quark masses and hence still obtain in turn somewhat too large a pion mass. The rapid improvement of computer power and performance supports the expectation that in the next years, fully realistic calculations will become feasible. – A more serious problem is that the method is at present restricted to studies at vanishing baryon number density. The reason for this is “purely technical”: for non-vanishing baryochemical potential $\mu$, the weight $\exp\{-S(U,\mu)\}$ is no longer positive definite, so that the standard Monte Carlo methods for generating equilibrium configurations breaks down. It is to be hoped that an alternative method will be found one of these years. This would then in particular allow us to address the interesting question of the phase structure of strongly interacting matter at high density and low temperature (see Fig. 1).

What have we learned so far from the computer simulation of finite temperature lattice QCD? The first observable to consider is the deconfinement measure $[5,6]$

$$L(T) \sim \exp\{-V(\infty)/T\}$$

(12)

where $V(r)$ is the potential between a static quark-antiquark pair separated by a distance $r$. In the limit of infinite current quark mass, i.e., in pure $SU(3)$ gauge theory, $L(T)$ becomes the order parameter of a center $Z_3$ symmetry, since $V(\infty) = \infty$. In the confinement regime, we therefore have $L = 0$; colour screening, on the other hand, makes $V(r)$ finite at large $r$, so that in the deconfined phase, $L$ does not vanish. In the large quark mass limit, deconfinement thus corresponds to the spontaneous breaking of a $Z_3$ symmetry, much like the onset of spontaneous magnetisation in a spin model. The structure of the theory here becomes that of a three-state Potts’ model, which shows a first order phase transition. One may thus expect a similar behaviour for $SU(3)$ gauge theory; for continuous transitions, one would say that spin and gauge theories are in the same universality class, and for $SU(2)$ gauge theory and the corresponding $Z_2$ spin system, the Ising model, this can in fact be shown. For finite current quark mass $m_q$, $V(r)$ remains finite for $r \to \infty$, since the string between the two colour charges “breaks” when the corresponding potential energy becomes equal to the mass $m_h$ of the lowest hadron; beyond this point, it becomes energetically more favourable to produce an additional hadron. Hence now $L$ no longer vanishes in the confined phase, but only becomes exponentially small there,

$$L(T) \sim \exp\{-m_h/T\};$$

(13)

here $m_h$ is typically of the order of the $\rho$-mass, since the pion as largely Goldstone boson plays a special role. This gives us $L \sim 10^{-2}$, rather than zero. Deconfinement is thus indeed much like the insulator-conductor transition, for which the order parameter, the
conductivity $\sigma(T)$, also does not really vanish for $T > 0$, but with $\sigma(T) \sim \exp\{-\Delta E/T\}$ is only exponentially small, since thermal ionisation (with ionisation energy $\Delta E$) produces a small number of unbound electrons even in the insulator phase.

Fig. 2 shows recent lattice results for $L(T)$ and the corresponding susceptibility $\chi_L(T) = \langle L^2 \rangle - \langle |L| \rangle^2$. The calculations were performed for the case of two flavours of light quarks, using a current quark mass about four times larger than that needed for the physical pion mass [7]. We note that $L(T)$ undergoes the expected sudden increase from a small confinement to a much larger deconfinement value. The sharp peak of $\chi_L(T)$ defines quite well the transition temperature; gauging the lattice scale in terms of the $\rho$-mass, we get $T_c \simeq 0.15$ GeV.

The next quantity to consider is the effective quark mass; it is measured by the expectation value of the corresponding term in the Lagrangian, $\langle \bar{\psi}\psi \rangle$. In the limit of vanishing current quark mass, the Lagrangian becomes chirally symmetric and $\langle \bar{\psi}\psi \rangle(T)$ the corresponding order parameter. In the confined phase, with effective constituent quark masses $m^{\text{const}}_q \simeq 0.3$ GeV, this chiral symmetry is spontaneously broken, while in the deconfined phase, at high enough temperature, we expect its restoration. In the real world, with finite pion and hence finite current quark mass, this symmetry is also only approximate, since $\langle \bar{\psi}\psi \rangle(T)$ now never vanishes at finite $T$.

The behaviour of $\langle \bar{\psi}\psi \rangle(T)$ and the corresponding susceptibility $\chi_m \sim \partial\langle \bar{\psi}\psi \rangle/\partial m_q$ are shown in Fig. 3, calculated for the same case as above in Fig. 2. We note here the expected sudden drop of the effective quark mass and the associated sharp peak in the susceptibility. The temperature at which this occurs coincides with that obtained through the deconfinement measure. We therefore conclude that at vanishing baryon number density, quark deconfinement and the shift from constituent to current quark mass coincide.

We thus obtain for $\mu = 0$ a rather well defined phase structure, consisting of a confined phase for $T < T_c$, with $L(T) \simeq 0$ and $\langle \bar{\psi}\psi \rangle \neq 0$, and a deconfined phase for $T > T_c$ with $L(T) \neq 0$ and $\langle \bar{\psi}\psi \rangle \simeq 0$. The underlying symmetries associated to the critical behaviour at $T = T_c$, the $Z_3$ symmetry of deconfinement and the chiral symmetry of the quark mass shift, become exact in the limits $m_q \to \infty$ and $m_q \to 0$, respectively. In the real world, both symmetries become approximate; nevertheless, we see from Figs. 2 and 3, that both associated measures retain an almost critical behaviour.

Next we turn to the behaviour of energy density $\epsilon$ and pressure $P$ at deconfinement. The most accurate results exist so far for pure gauge theory, i.e., for QCD in the infinite quark mass limit; the qualitative behaviour remains much the same when quarks are included. We see in Fig. 4 that $\epsilon/T^4$ in SU(3) gauge theory [8] changes quite abruptly at the above determined critical temperature $T = T_c$, increasing from a low hadronic value (difficult to determine precisely with the accuracy of present lattice calculations) to nearly that expected for an ideal gas of quarks and gluons. The forms shown here are obtained by extrapolating results from different lattice size calculations to the continuum limit. Besides the sudden increase at deconfinement, there are two further points to note. The interaction measure $(\epsilon - 3P)/T^4$, if correctly normalized, vanishes for an ideal gas of massless particles. It certainly does not vanish here, as shown more explicitly in Fig. 5. In particular, it is quite strongly peaked in the region $T_c < T < 2 T_c$, indicating
the presence of strong remaining interaction effects in this region. The nature of these effects is presently one of the main questions in finite temperature lattice QCD; there are a number of model proposals to account for the “measured” lattice data. The second, somewhat surprising effect is that the thermodynamic observables do not fully attain their Stefan-Boltzmann values (marked “SB” in Fig. 4) even at very high temperatures, in contrast to earlier conclusions based on less precise calculations. The remaining 10 to 15% deviation could well be due to effective “thermal” masses of gluons (and of quarks in full QCD); this problem also is in under investigation.

Finally we turn to the value of the transition temperature. Since QCD (in the limit of massless quarks) does not contain any dimensional parameters, \( T_c \) can only be obtained in physical units by expressing it in terms of some other known observable which can also be calculated on the lattice, such as the \( \rho \)-mass, the proton mass, or the string tension. In the continuum limit, all different ways should lead to the same result. Within the present accuracy, they define the uncertainty so far still inherent in the lattice evaluation of QCD. Using the \( \rho \)-mass to fix the scale leads to \( T_c \simeq 0.15 \text{ GeV} \), while the string tension still allows values as large as \( T_c \simeq 0.20 \text{ GeV} \). This means that energy densities of some 1 - 2 GeV/fm\(^3\) are needed in order to produce a medium of deconfined quarks and gluons.

In summary, finite temperature lattice QCD shows

- that there is a deconfinement transition with an associated shift in the effective quark mass at \( T_c \simeq 0.15 - 0.20 \text{ GeV} \);
- this transition is accompanied by a sudden increase in the energy density (“latent heat of deconfinement”) from a small hadronic value to a much larger value some ten to twenty percent below that of ideal quark-gluon plasma;
- for \( T_c \leq T \leq 2 T_c \), the ideal gas measure \( (\varepsilon - 3P)/T^4 \) differs very much from zero, indicating the presence of considerable plasma interactions.

Both conceptually and through \textit{ab initio} QCD calculations we thus have a relatively good understanding of the critical behaviour and the phase structure expected for strongly interacting matter.

3. Colour Deconfinement in Nuclear Collisions

How can these predictions be tested? Since ten years, experiments studying the collision of nuclei at high energies are carried out with the aim of producing strongly interacting matter in the laboratory. There are today promising indications that the hadrons produced in such collisions indeed come from equilibrated thermal systems. In the next section, we want to address in particular the question of how one can check if these systems in their early hot and dense stages consisted of deconfined quarks and gluons. In case of thermal evolution, the hadrons formed in the final stage do not carry any information about the earlier stages of the system. We thus need to find a probe which somehow retains this primordial information. The most promising and most intensively studied probe of this kind is the behaviour of charmonium production in nuclear collisions, proposed ten years ago [9]. With the help of this probe, very recent experimental results may provide the first indication of the onset of colour deconfinement.
A hadron placed into a deconfining medium will dissolve into its quark constituents. If the medium expands, cools off and eventually hadronizes, normal hadrons will now reappear and, in case of an expansion in thermal equilibrium, not carry any information about the earlier stages.

The situation is quite different for a $J/\psi$ put into a quark-gluon plasma. The $J/\psi$ is a bound state of the heavy $c$ and $\bar{c}$ quark, which each have a mass of about 1.5 GeV. The $J$ has a mass of about 3.1 GeV; with a radius of about 0.3 fm it is much smaller than the usual light hadrons. It has a binding energy (the difference between $J/\psi$ mass and open charm threshold) of about 0.64 GeV, which is much larger than the typical hadronic scale $\Lambda_{\text{QCD}} \simeq 0.2$ GeV. The $J/\psi$ is produced quite rarely in hadronic collisions – at present energies, in about one out of $10^5$ events. Through its decay into dimuons, it is, however, rather easily detectable in suitably triggered experiments, so that in the ongoing heavy ion studies at CERN [10], some hundred thousand $J/\psi$’s are measured for a given target-projectile combination.

If the quark-gluon plasma is sufficiently hot, a $J/\psi$ will also melt in it. However, its constituents, the $c$ and the $\bar{c}$, now separate and never meet again. Since the production of more than one $c\bar{c}$ pair per collision is very strongly excluded, the $c$ must at hadronisation combine with a normal antiquark, the $\bar{c}$ with a normal quark, leading to a $D$ and a $\bar{D}$, respectively. If nuclear collisions produce a deconfining medium, then such collisions must also lead to a suppression of $J/\psi$ production [9].

Before we can use this as a deconfinement probe, we must know if there are possible interactions in a confined medium which can cause $J/\psi$ suppression. This question is now answered, including in particular also a direct experimental test of the theoretical answer. The cross section for the dissociation of a $J/\psi$ colliding with a usual light hadron can be calculated in short distance QCD [11]; the relevant diagram is shown in Fig. 6.

Fig. 6: $J/\psi$ interaction with a light hadron
It contains two parts: the break-up of a $J/\psi$ by an incident gluon, essentially the photo-effect analog in QCD, and the emission or absorption of a gluon by a light hadron. The latter is evidently of non-perturbative nature, but the needed gluon distribution function is determined in deep inelastic scattering. Starting from the operator product expansion, one thus establishes sum rules relating the dissociation cross section $\sigma_{h-\psi}^{\text{in}}$ to the gluon distribution function $g_h(x)$ for a light hadron. The large binding energy of the $J/\psi$ requires hard gluons for both resolution and break-up. On the other hand, the presence of hard gluons in light hadrons of present momenta is strongly suppressed; $g_h(x)$ falls rapidly for large $x$. As a result, $\sigma_{h-\psi}^{\text{in}}(s)$ suffers a very strong threshold damping; it is essentially zero until the $h-\psi$ collision energy $\sqrt{s}$ becomes much larger than the presently available 20 GeV (Fig. 7).

We thus conclude that in the presently produced media, collisions with hadrons cannot dissociate $J/\psi$’s. This makes $J/\psi$ suppression into an unambiguous deconfinement probe: $J/\psi$’s can be suppressed in a given medium if and only if this medium contains deconfined gluons.

Since the predicted threshold suppression of $\sigma_{h-\psi}^{\text{in}}(s)$ is crucial in this argument, it should certainly be tested experimentally. So far, it has been confirmed for the very similar process of $J/\psi$ photo-production. A direct check is possible, however, through an “inverse kinematics” experiment, shooting a heavy nuclear beam at a hydrogen or deuterium target [12]. Such an experiment should certainly be carried out, and preparations are in progress.

Before we can actually analyse the confinement status of the media formed in nuclear collisions, one further problem has to be addressed. We have so far implicitly considered the fate of a fully formed physical $J/\psi$ in the given medium. However, the same collision that produces the medium also has to produce the $J/\psi$, and neither production process is instantaneous. The mechanism of $J/\psi$ production in hadronic collisions has recently been established more precisely, both experimentally [13] and theoretically [14,15]. The first stage of the production process is the formation of a coloured $c\bar{c}$ pair, which combines with a colliner gluon to form a colour-neutral $c\bar{c} - g$ state. This pre-resonance charmonium state turns into a physical $J/\psi$, i.e., a colour singlet $c\bar{c}$ state, after some 0.2 to 0.3 fm. In its pre-resonance stage, however, it can interact with nuclear matter, and this interaction must be taken into account before studying the suppression of physical $J/\psi$’s. Theoretical estimates [15] give the $c\bar{c} - g$ state in interactions with hadrons a dissociation cross section of some 6 - 7 mb. This can be studied in $p - A$ interactions, and recent precision data [16,17] confirm the picture both qualitatively and quantitatively.

The survival probability of the pre-resonance $c\bar{c} - g$ state passing through a length $L$ of normal nuclear matter can be estimated by $S_{J/\psi}(L) = \exp\{-n_0 \sigma_{c\bar{c}g} L\}$; this estimate can be checked by calculations based on Glauber theory. Looking at the survival probability in $p - A$ collisions as function of $L \simeq (3/4) R_A^{1/3}$ leads to a cross section $\sigma_{c\bar{c}g} = 6.3 \pm 1.2$ mb for the break-up of pre-resonance charmonia by collisions with nucleons (Fig. 8), which agrees with the theoretical expectations.

We thus find that $J/\psi$ suppression in $p - A$ collisions is well described in terms of pre-resonance dissociation. The equality of $J/\psi$ and $\psi'$ suppression in such interactions
moreover finds a natural explanation by such a process. Turning now to nuclear collisions and the question of deconfinement, we have to check if such reactions lead to a suppression beyond the known $c\bar{c} - g$ dissociation in normal nuclear matter. The relevant data from $O - Cu$, $O - U$ and $S - U$ interactions at the CERN-SPS are included in Fig. 8; we see that they agree completely with the predicted pre-resonance suppression. This result can also be extended to $S - U$ collisions at different impact parameters. Hence we conclude that up to central $S - U$ collisions, the $J/\psi$ does not suffer any nuclear effect beyond the mentioned pre-resonance suppression in nuclear matter. In other words, these collisions do not produce a deconfined medium, even though the associated average energy density is expected to be in the range of $1 - 3$ GeV/fm$^3$.

A quark-gluon plasma must thus lead to more $J/\psi$ suppression than that obtained from the pre-resonance $c\bar{c} - g$ absorption. The recent announcement [16,17] of a strong “anomalous” $J/\psi$ suppression in $Pb-Pb$ collisions at CERN may therefore provide a first hint of deconfinement. Although the average energy density in $Pb-Pb$ collisions is only slightly higher than that in central $S-U$ interactions, the NA50 collaboration finds, as function of collision centrality, a very rapid onset of much stronger suppression. In particular, central $Pb-Pb$ collisions result in a suppression which is more than twice that due to pre-resonance absorption in normal nuclear matter, as shown in Fig. 9. The path length $L$ used as a variable there provides in nucleus-nucleus collisions also a measure of the average energy density $\epsilon$; hence the anomalous suppression sets in suddenly at a certain value of $\epsilon$.

A simple model can illustrate how such an effect could arise [18,19]. Although the average energy densities in central $S-U$ and central $Pb-Pb$ collisions are very similar, the corresponding energy density profiles are quite different (Fig. 10); the interior of the interaction region is much hotter in $Pb-Pb$ than in $S-U$ collisions. If the peak value in central $S-U$ collisions is just the critical energy density needed for $J/\psi$ melting, then all $J/\psi$’s formed in the $Pb-Pb$ region hotter than this will be suppressed. It turns out that this in fact gives just the right amount of anomalous $J/\psi$ suppression [19].

Clearly more studies, both experimental and theoretical, are needed before a definite conclusion is possible. At this time, however, it does seem that if one can confirm

- the sudden onset of the anomalous suppression, e.g. by more peripheral $Pb-Pb$ collisions, different $A-B$ combinations, different collision energies, and
- the transparency of confined matter to $J/\psi$’s by an inverse kinematics experiment, then the observed effect would seem to be the first indication of colour deconfinement in nuclear collisions.

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$5.20 \sim 5.30$

$0 \sim 0.1 \sim 0.2 \sim 0.3$

$m_q a = 0.02$

$\chi_L$

$L$

$6/g^2 \sim T$
$m_q a = 0.02$
\((\varepsilon - 3p)/T^4\)
\[ J/\psi \text{ cross-sections versus } L \]

\[ \frac{B_{\mu\mu}\sigma(J/\psi)}{A_{\text{proj.}}B_{\text{target}}} \left( \text{nb/nucleon}^2 \right) \sim \exp\left( -\rho L \sigma_{\text{abs}} \right) \]

\[ \rho = 0.138 \text{ n/fm}^3 \]

\[ \sigma_{\text{abs}} = 6.3 \pm 1.2 \text{ mb} \]

*rescaled to 200 GeV/c*
$\epsilon$ [GeV/fm$^3$]

- $\text{Pb-Pb}$ region
- $\text{S-U}$ region

$r$ [fm]

$\epsilon$ [GeV/fm$^3$]