Colour Deconfinement and $J/\psi$ Suppression in High Energy Nuclear Collisions* 

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1. Introduction 

Strong interaction thermodynamics deals with the behaviour of matter at extreme temperatures and densities. Its central theme is the transition from hadronic matter to a new state, the quark-gluon plasma, in which quarks are no longer confined to colour-neutral bound states. The existence of the quark-gluon plasma (QGP) is one of the basic predictions of quantum chromodynamics (QCD), and its experimental observation represents one of the great challenges to present high energy physics.  

Colour deconfinement is essentially the QCD version of the insulator-conductor transition familiar from condensed matter physics: at sufficiently high colour charge densities, screening suppresses the long-range confining part of the strong interaction, 

$$\sigma r \to \sigma r \left( 1 - \frac{e^{-\mu r}}{\mu r} \right),$$  \hspace{1cm} (1) 

where $\mu^{-1}$ is the screening radius. Such screening leads to the melting of hadronic states, as illustrated in Fig. 1. 

*) Lecture given at the 35th Course of the International School of Subnuclear Physics, Erice/Sicily, 26. 8. – 4. 9. 1997, to appear in the Proceedings.
Figure 1: Deconfinement by colour screening.

The computer simulation of finite temperature QCD on the lattice allows an \textit{ab initio} study of this phenomenon, starting from the QCD Lagrangian as dynamical basis. We cite here only the most important results of such studies \cite{1}. As seen in Fig. 2, the energy density of strongly interacting matter shows a sudden increase at a temperature \( T_c \simeq 150 \text{ MeV} \); from a value near that for an ideal gas of pions, it jumps to that for an ideal plasma of quarks and gluons \cite{2}. The associated pressure, however, follows this increase more slowly; since \( \epsilon = 3P \) for an ideal gas of massless constituents, this indicates that at least until \( T/T_c \simeq 2 - 3 \) there must still be considerable remnant interactions in the plasma. In Fig. 3, we see that the increase of the energy density is indeed due to deconfinement. The order parameter for quark binding is 
\[ \langle L \rangle \sim \exp\left\{-\frac{V(\infty)}{T}\right\}, \]
where \( V(\infty) \) is the potential of a static quark-antiquark pair in the limit of infinite separation. In confined matter, \( V(\infty) \) diverges, causing \( \langle L \rangle \) to vanish; in a QGP, screening keeps \( V(\infty) \) and hence also \( \langle L \rangle \) finite. We see that the transition for \( \langle L \rangle \) coincides with the increase of \( \epsilon \). Moreover, we note that at the same temperature the effective quark mass, measured by \( \langle \bar{\psi} \psi \rangle \), drops from the finite value associated with constituent quarks to the essentially vanishing value of the light \( u \) and \( d \) quarks. From this, we can conclude that in the regime studied here (vanishing overall baryon number density), deconfinement and chiral symmetry restoration coincide.

Lattice QCD thus predicts that the transition from hadronic matter to a QGP occurs at a temperature of around 150 - 200 MeV. The uncertainty in this value is due to finite lattice sizes and to uncertainties in the calculation of hadron masses; both can be removed in the next few years by improved computer technology. In any case, we know already now (see Fig. 2) that in order to produce a quark-gluon plasma, energy densities of about 1 - 3 GeV/fm\(^3\) are necessary. Shortly after the big bang, our universe must have consisted of such matter, and it would certainly be interesting to reproduce it in the laboratory.
Figure 2: Energy density and pressure in two-flavour QCD [2].

Figure 3: Deconfinement and chiral symmetry in two-flavour QCD [3].
High energy nuclear collisions are expected to provide the required conditions. When two nuclei collide at high energy, they pass through each other and leave behind a trail of locally disturbed vacuum, of energy deposited in bubbles of strongly interacting matter; these bubbles are our candidates for the ‘little bang’. They were formed shortly after the passage of two nuclei, and the first question is whether they were sufficiently hot and dense to reach colour deconfinement. In any case, since they are not contained by anything, they will expand, cool off and eventually freeze out into free hadrons, which can be observed experimentally. We can use these hadrons to estimate the initial energy density of the bubbles [3], by letting the evolution film run backwards. The observed number \( (dN/dy) \) of produced hadrons, each of average energy \( p_0 \), must have originated in a volume of transverse size determined by the colliding nuclei; the longitudinal size can be estimated if the bubbles are defined to give isotropic momentum distributions. In this way, we obtain for an \( A - A \) collision

\[
\epsilon_0 \simeq \left( \frac{dN}{dy} \right)_{y=0} \frac{p_0}{\pi R_A^2 \tau_0},
\]

where \( R_A \simeq 1.12A^{1/3} \) and \( \tau_0 \simeq 1 \text{ fm} \). Using the multiplicites measured at CERN-SPS energies, we get energy densities in the range 2 - 5 GeV/fm\(^3\); estimates for RHIC and LHC run as high as 10 - 20 GeV/fm\(^3\). Thus there seems to be a real chance to produce and study the QGP in the laboratory, making the quark-hadron transition the only experimentally accessible cosmological phase transition.

The fundamental question facing us is therefore how to check if the medium produced in high energy nuclear collisions was deconfined in its early stages. To show how non-trivial this problem is, we recall that no measurement on a glass of water can ever tell us if it was ice half an hour ago. We want to find out what things were like before the time at which we can carry out measurements. A useful deconfinement probe thus has to fulfill a number of conditions:

- It must be present early and retain its memory throughout the evolution;
- it must be hard enough to resolve the short-distance scales of the QGP; and
- it must be able to distinguish confined and deconfined media.

So far, two candidates have been proposed to do this:

- quarkonium states, whose dissociation pattern is quite different in confined and deconfined media (“\( J/\psi \) suppression” [4]), and
- hard jets, whose energy loss and transverse momentum broadening depends on the nature of the medium they traverse (“jet quenching” [5]).

We shall here consider the first of these two and study quarkonium dissociation in nuclear collisions as test for colour deconfinement. It gives rather clear-cut theoretical predictions for experiments accessible at SPS energies; jets will probably have to wait for RHIC or LHC studies.
Before turning to the details of the probe, we restate the basic question: Given a box of unidentified matter, we want to know if the quarks and gluons which make up this matter are in a confined or deconfined state. It is thus not evidence for quarks and gluons that we look for, but rather the confinement/deconfinement status of these elementary building blocks of all forms of matter.

2. Charmonium Dissociation and Colour Deconfinement

As prototype for charmonium, consider the $J/\psi$; it is the $1S$ bound state of a charm quark ($m_c \simeq 1.4$ GeV) and its antiquark, with $M_{J/\psi} \simeq 3.1$ GeV. Its usefulness as deconfinement probe is easily seen. If a $J/\psi$ is placed into a hot medium of deconfined quarks and gluons, colour screening will dissolve the binding, so that the $c$ and the $\bar{c}$ separate. When the medium cools down to the confinement transition point, they will therefore in general be too far apart to see each other. Since thermal production of further $c\bar{c}$ pairs is negligibly small because of the high charm quark mass, the $c$ must combine with a light antiquark to form a $D$, and the $\bar{c}$ with a light quark for a $\bar{D}$. The presence of a quark-gluon plasma will thus lead to a suppression of $J/\psi$ production [4].

We shall now first consider this $J/\psi$ suppression in terms of colour screening, i.e., as consequence of global features of the medium, and then turn to a microscopic approach, in which the bound state is assumed to be broken up by collisions with constituents of the medium.

2a. Colour Screening

Because of the large charm quark mass, the charmonium spectrum can be calculated with good precision by means of the Schrödinger equation [6]

$$[2m_c + \frac{1}{m_c} \nabla^2 + V(r)]\Psi_{n,l} = M_{n,l}\Psi_{n,l},$$

(3)

where the potential $V(r) = \sigma r - \alpha/r$ contains a confining long-distance part $\sigma r$ and a Coulomb-like short-distance term $\alpha/r$. For different values of the principal quantum number $n$ and the orbital quantum number $l$, the masses $M_{n,l}$ and the wave functions $\Psi_{n,l}(r)$ of different charmonium states $J/\psi$, $\chi$, $\psi'$, ... in vacuum are given in terms of the constants $m_c$, $\sigma$ and $\alpha$.

In a medium, the potential becomes screened,

$$V(r, \mu) = \frac{\sigma}{\mu}[1 - e^{-\mu r}] - \frac{\alpha}{r}e^{-\mu r},$$

(4)

where $\mu$ is the screening mass, i.e., $r_D = \mu^{-1}$ is the ‘Debye’ colour screening radius. Screening is a global feature of the medium, shortening the range of the binding potential. Once $\mu$ becomes sufficiently large, the bound states begin to disappear, starting with the most weakly bound; hence for $\mu \geq \mu_d$, the bound state $i$ is no longer possible [7].

Using finite temperature lattice QCD, we can now determine the screening mass $\mu(T)$ as function of the temperature $T$ or, equivalently, as function of the energy density $\epsilon$ of the medium [8]. The resulting melting pattern for the most important charmonium states is
summarized in Table 1; we see that while both $\psi'$ and $\chi$ melt at the critical deconfinement point, the $J/\psi$, being smaller and more tightly bound, survives to about $1.2T_c$ and hence about twice the critical energy density. With increasing temperature, a hot medium will thus lead to successive charmonium melting, so that the suppression or survival of specific charmonium states serves as a thermometer for the medium, in much the same way as the relative intensity of spectral lines in stellar interiors measure the temperature of stellar matter \[9\]. Note, however, that other possible sources of charmonium dissociation have to be considered before the method becomes unambiguous. This will be done in the next subsection.

| state | $\mu_d$ [GeV] | $T_d$ | $\epsilon_d$ |
|-------|---------------|-------|-------------|
| $J/\psi$ | 0.70 | $1.2T_c$ | $2\epsilon_c$ |
| $\chi$ | 0.35 | $T_c$ | $\epsilon_c$ |
| $\psi'$ | 0.35 | $T_c$ | $\epsilon_c$ |

Table 1: Charmonium Dissociation by Colour Screening

2b. Gluon Dissociation

The binding energy of the $J/\psi$, i.e., the energy difference between the $J/\psi$ mass and the open charm threshold, is with $\Delta E_{J/\psi} = 2M_D - M_{J/\psi} \simeq 0.64$ GeV $\gg \Lambda_{QCD}$ considerably larger than the typical non-perturbative hadronic scale $\Lambda_{QCD} \simeq 0.2$ GeV. As a consequence, the size of the $J/\psi$ is much smaller than that of a typical hadron, $r_{J/\psi} \simeq 0.2$ fm $\ll \Lambda_{QCD}^{-1} = 1$ fm. Hence the $J/\psi$ is a hadron with characteristic short-distance features; in particular, rather hard gluons are necessary to resolve or dissociate it, making such a dissociation accessible to perturbative calculations. $J/\psi$ collisions with ordinary hadrons made up of the usual $u, d$ and $s$ quarks thus probe the local partonic structure of these ‘light’ hadrons, not their global hadronic aspects, such as mass or size. It is for this reason that $J/\psi$’s can be used as confinement/deconfinement probe.

This can be illustrated by a simple example. Consider an ideal pion gas as a confined medium. The momentum spectrum of pions has the Boltzmann form $f(p) \sim \exp(-(|p|/T))$, giving the pions an average momentum $\langle |p| \rangle = 3T$. With the pionic gluon distribution function $xg(x) \sim (1 - x)^3$, where $x = k/p$ denotes the fraction of the pion momentum carried by a gluon, the average momenta of gluons confined to pions becomes

$$\langle |k| \rangle_{\text{conf}} \simeq 0.6T.$$  \hspace{1cm} (5)

On the other hand, an ideal QGP as prototype of a deconfined medium gives the gluons themselves the Boltzmann distribution $f(k) \sim \exp(-(|k|/T))$ and hence average momenta

$$\langle |k| \rangle_{\text{deconf}} = 3T.$$  \hspace{1cm} (6)

Deconfinement thus results in a hardening of the gluon momentum distribution. More generally speaking, the onset of deconfinement will lead to parton distribution functions
which are different from those for free hadrons, as determined by deep inelastic scattering experiments. Since hard gluons are needed to resolve and dissociate \( J/\psi \)'s, one can use \( J/\psi \)'s to probe the in-medium gluon hardness and hence the confinement status of the medium.

These qualitative considerations can be put on a solid theoretical basis provided by short-distance QCD \([10] - [13]\). In Fig. 4 we show the relevant diagram for the calculation of the inelastic \( J/\psi \)-hadron cross section, as obtained in the operator product expansion (essentially a multipole expansion for the charmonium quark-antiquark system). The upper part of the figure shows \( J/\psi \) dissociation by gluon interaction; the cross section for this process,

\[
\sigma_{g-J/\psi} \sim (k - \Delta E_{J/\psi})^{3/2}k^{-5},
\]

constitutes the QCD analogue of the photo-effect. Convoluting the \( J/\psi \) gluon-dissociation with the gluon distribution in the incident hadron, \( xg(x) \approx 0.5(1-x)^n \), we obtain

\[
\sigma_{h-J/\psi} \approx \sigma_{\text{geom}}(1 - \lambda_0/\lambda)^{n+2.5}
\]

for the inelastic \( J/\psi \)-hadron cross section, with \( \lambda \approx (s - M^2_{\psi})/M_{\psi} \) and \( \lambda_0 \approx (M_h + \Delta E_{\psi}) \); \( s \) denotes the squared \( J/\psi \)-hadron collision energy. In Eq. (8), \( \sigma_{\text{geom}} \approx \text{const.} \). In Eq. (8), we obtain \( \sigma_{h-J/\psi} \approx \sigma_{\text{geom}}(1 - \lambda_0/\lambda)^{n+2.5} \) in Eq. (8). In Fig. 5, we compare the cross sections for \( J/\psi \) dissociation by gluons (“gluo-effect”) and by pions \((n = 3)\), as given by Eq’s (7) and (8). Gluon dissociation shows the typical photo-effect form, vanishing until the gluon momentum \( k \) passes the binding energy \( \Delta E_{\psi} \); it peaks just a little later and then vanishes again when sufficiently hard gluons just pass through the much larger charmonium bound states. In contrast, the \( J/\psi \)-hadron cross section remains negligibly small until rather high hadron momenta \((3 - 4 \text{ GeV})\). In a thermal medium, such momenta correspond to temperatures of more than one GeV. Hence confined media in the temperature range of a few hundred MeV are essentially transparent to \( J/\psi \)'s, while deconfined media of the same temperatures very effectively dissociate them and thus are \( J/\psi \)-opaque.

\[\text{Figure 4: } J/\psi \text{ dissociation by hadron interaction.}\]
The situation for $\chi$’s is quite similar, except for an earlier gluon dissociation threshold due to the lower $\chi$ binding energy of about 250 MeV. This difference provides the microscopic basis for the successive melting of different charmonium states noted above. – For the $\psi'$, however, we have an almost negligible binding energy of only 60 MeV, so that there cannot really be any difference in its dissociation by confined or deconfined media. In other words, any strongly interacting matter is expected to be $\psi'$-opaque.

![Graph showing J/ψ dissociation by gluons and by pions](image)

**Figure 5:** $J/\psi$ dissociation by gluons and by pions; $k$ denotes the momentum of the projectile incident on a stationary $J/\psi$.

We can thus define a schematic test to determine the confinement status of a box of unidentified matter. We first shine a $\psi'$ beam at it: if it is transparent to this, the box does not contain strongly interacting matter; otherwise it does. In that case we repeat the procedure with a $J/\psi$ beam: if this is unaffected, the medium in the box is confined; if the beam is attenuated, it is deconfined. The problem of an unambiguous deconfinement test is thus in principle solved: there is $J/\psi$ dissociation if and only if the medium is deconfined. However, the test pre-supposes the existence of a prepared strongly interacting medium and the availability of $J/\psi$ and $\psi'$ beams as probes, and in
nuclear collisions, neither of these is immediately given.

3. $J/\psi$ Production in Nuclear Collisions

In nuclear collisions, both the charmonium states and the matter to be probed are produced in the course of the collision, and both require a finite ‘formation time’ to be formed. We therefore have to ask what will happen ‘before’ they are there and consider in particular pre-resonance absorption in normal nuclear matter.

Quarkonium production in hadron-hadron collisions has in recent years been studied quite extensively, triggered in particular by detailed and quite conclusive experiments at FNAL [14]. The production of a $J/\psi$ in a $p−p$ or $p−\bar{p}$ collision begins with the production of a $c\bar{c}$ pair, which occurs at high energies dominantly by gluon fusion (Fig. 6). The $c\bar{c}$ is generally in a colour octet state and has to neutralize its colour in order to leave the interaction region and form a $J/\psi$. The colour singlet model [15] proposed a perturbative treatment of this colour neutralisation; this is now clearly ruled out by the mentioned FNAL experiments. Colour neutralisation thus takes place in a non-perturbative way [16], and the recently proposed colour octet formalism [17] can be extended to obtain such a description [18]. The coloured $c\bar{c}$ binds with a soft collinear gluon to form a colour singlet $c\bar{c}−g$ state; in a proper time $\tau_{c\bar{c}g} \approx (2m_c\Lambda_{\text{QCD}})^{-1/2} \approx 0.3$ fm, this falls into the $c\bar{c}$ singlet state which is the basic component of a physical $J/\psi$. Since the life-time (and hence also the size) of this pre-resonance $c\bar{c}g$ state is determined by the hadronic scale $\Lambda_{\text{QCD}}$, it is the same for $J/\psi$, $\chi$ and $\psi'$; in contrast, the final resonances have very different geometric sizes.

Figure 6: $J/\psi$ production by $p−p$ collision.

This formation process has rather striking consequences for charmonium production in $p−A$ collisions. The presence of the nuclear target medium is known to reduce $J/\psi$ production rates in $p−A$ collisions [19, 20], relative to those in $p−p$ interactions (Fig. 7). However, these experiments are carried out in a kinematic region giving the
nascent $J/\psi$ momenta of 50 GeV/c or more in the target rest frame. As a result, the transition $c\bar{c}g \rightarrow J/\psi$, $\chi$ or $\psi'$ occurs outside the target nucleus; the nuclear matter of the target sees only the passage of the $c\bar{c}g$ state. Hence the observed attenuation of charmonium production should be the same for $J/\psi$ as for $\psi'$, as it is indeed found to be (see Fig. 8 [21]). Earlier attempts to explain charmonium suppression in $p - A$ interactions in terms of the absorption of physical $J/\psi$ states [22] had encountered difficulties precisely because of this feature. The equal attenuation of $J/\psi$ and $\psi'$ is a natural consequence of pre-resonance absorption; it can never be obtained for the physical $J/\psi$ and $\psi'$ states with their very different geometric sizes.

The cross section for the dissociation of the $c\bar{c}g$ state through collisions in nuclear matter can be estimated theoretically [18]; but it can also be determined directly from $p - A$ data [23]. A $c\bar{c}$ pair formed at point $z_0$ in the target nucleus has a survival probability

$$S_{c\bar{c}g}^A = \exp - \left\{ \int_{z_0}^{\infty} dz \rho_A(z) \sigma_{c\bar{c}g - N} \right\},$$

where the integration covers the path remaining from $z_0$ out of the nucleus. The traversed medium of nucleus $A$ is parametrized through a Wood-Saxon density distribution $\rho_A(z)$, and by comparing $S_{c\bar{c}g}^A$ with data for different targets $A$, the dissociation cross section for $c\bar{c}g - N$ interactions is found to be [23]

$$\sigma_{c\bar{c}g - N} = 7.3 \pm 0.6 \text{ mb}.$$  (10)

In Fig. 9 it is seen that pre-resonance absorption with this cross section agrees well with all presently available $p - A$ data on $J/\psi$ production.
Figure 8: The relative $A$-dependence of $J/\psi$ and $\psi'$ production.

Figure 9: $J/\psi$ production in $p - A$ collisions, compared to pre-resonance absorption in nuclear matter [23].
We now turn to nucleus-nucleus collisions; here there will certainly also be pre-resonance absorption in nuclear matter. However, in addition to the target and projectile nuclei, there could now be a substantial amount of produced ‘secondary’ medium (Fig. 10). We want to check if there is such a medium and if yes, test its confinement status.

Figure 10: Schematic view of $J/\psi$ production in $A-B$ collisions.

The survival probability of a pre-resonance charmonium state in an $A-B$ collision at impact parameter $b$ is given by

$$S_{c\bar{c}g}^{AB}(b) = \exp \left\{ - \int_{z_0^A}^{\infty} dz \rho_A(z) \sigma_{c\bar{c}g-N} + \int_{z_0^B}^{\infty} dz \rho_B(z) \sigma_{c\bar{c}g-N} \right\},$$  \hspace{1cm} (11)

in extension of Eq. (9). Here $z_0^A$ specifies the formation point of the $c\bar{c}g$ within nucleus $A$, $z_0^B$ its position in $B$. Since experiments cannot directly measure the impact parameter $b$, we have to specify how Eq. (10) can be applied to data.

The Glauber formalism allows us to calculate the number $N_{w}^{AB}(b)$ of participant (‘wounded’) nucleons for a given collision. The number of secondary hadrons produced in association with the observed $J/\psi$ is found to be proportional to $N_{w}^{AB}$. The transverse energy $E_T$ carried by the secondaries is measured experimentally, together with the $J/\psi$’s. We thus have

$$E_T(b) = q N_{w}^{AB}(b);$$  \hspace{1cm} (12)

the proportionality constant $q$ has to be determined on the basis of the given experimental acceptance. Once it is fixed, we have to check that the collision geometry (the measured relation between $E_T$ and the number of spectator nucleons, the measured $E_T$-distribution) are correctly reproduced [23]. Once this is assured, we can check if the $J/\psi$ production in $O-Cu$, $O-U$ and $S-U$, as measured by the NA38 experiment at CERN over the past ten years [24], shows anything beyond the expected pre-resonance nuclear absorption with the cross section $\sigma_{c\bar{c}g-N} = 7.3 \pm 0.6$ mb determined from $p-A$ interactions.

The answer is clearly negative, as seen in Fig. 11 for the integrated cross sections and in Fig. 12 for the centrality ($E_T$) dependence of $S-U$ collisions. All nuclear collisions measured by NA38 thus show only what is now called ‘normal’ $J/\psi$ suppression, i.e., the pre-resonance suppression already observed in $p-A$ interactions [18, 22, 23]. We thus
have to ask if the $A - B$ collisions studied by NA38 lead to any produced secondary medium at all. The behaviour of $\psi'$ production (Fig. 13) shows that this is indeed the case [23]: there is $\psi'$ suppression beyond the expected pre-resonance nuclear absorption. There thus is a secondary medium, and it can distinguish a $\psi'$ (which is suppressed) from a $J/\psi$ (which remains unaffected). In other words, this medium shows no indication for colour deconfinement.

4. Anomalous $J/\psi$ Suppression

In view of this state of affairs, the 30% additional (and hence ‘anomalous’) $J/\psi$ suppression observed in $Pb - Pb$ collisions by the NA50 collaboration at CERN [26] caused considerable excitement; the results from the first (1995) run are shown in Fig’s. 14 and 15. Is this the first indication for the onset of colour deconfinement [27, 28]?

To address this question, we have to check if it is possible to interpret the observed effect in terms of dissociation in a confined medium. The $Pb - Pb$ collision could produce a medium of secondary hadrons, and the $J/\psi$ could be broken up by interactions with these hadronic comovers 1. The main feature of such absorption is that it has no specific onset, but is always operative; hence it would have to be present to some extent also in $S - U$ collisions. Its functional form at fixed impact parameter $b$ is [29, 30]

$$S_{co}(b) = \exp \{ - \{ n_{co}(b) \sigma_{co} \tau_0 \ln[n(\epsilon)/n_f] \} \}, \tag{13}$$

where $n(b)$ is the comover density, $\tau_0 \simeq 1$ fm the time required to form the comover medium, and $n_f$ the comover density at freeze-out. Eq. (13) always leads to a monotonic increase of suppression with increasing centrality and comover density.

$J/\psi$ dissociation by colour deconfinement, in contrast, sets in at a specific energy density of the deconfining medium (see Table 1 and Fig. 16). To illustrate the consequences, consider a central $A - A$ collision. The energy density in the transverse plane will be highest at the center, along the collision axis, and will decrease with increasing distance $r$ from the axis (Fig. 17). Once the central value passes the critical deconfinement point, the $J/\psi$'s in the ‘hot’ central region start melting, while those in the cooler outer rim survive. If we take a hard sphere nuclear distribution, this leads to

$$S_{\psi}(\epsilon) = \Theta(\epsilon - \epsilon_{\psi}) + \Theta(\epsilon_{\psi} - \epsilon) \left( \frac{\epsilon_{\psi}}{\epsilon} \right)^{9/4}, \tag{14}$$

for the survival probability of the $J/\psi$ in a medium of central energy density $\epsilon = \epsilon(r = 0)$ [31]. If we identify the deconfinement threshold $\epsilon_{\psi}$ with the energy density at $r = 0$ in a central $S - U$ collision and use more a realistic Wood-Saxon potential, then this reproduces quite well the amount of suppression observed in central $Pb - Pb$ collisions [27].

1 Such an approach is evidently not in accord with the small hadron-$J/\psi$ dissociation cross section obtained in Eq. (8); it thus implicitly assumes that the charm quark mass is not yet large enough for the applicability of short-distance QCD calculations.
Figure 11: $J/\psi$ production in $A - B$ collisions, compared to pre-resonance absorption in nuclear matter [23].

Figure 12: The $E_T$-dependence of $J/\psi$ production in $S - U$ collisions, compared to pre-resonance absorption in nuclear matter [23].
Figure 13: $J/\psi$ and $\psi'$ production in $A - B$ collisions [25], compared to pre-resonance absorption in nuclear matter [23].

To study the onset of deconfinement in more detail, we have to take into account the fact that of the observed $J/\psi$'s, about 60% are directly produced $1S$ states, while the remainder is due to $\chi$ ($\sim 30\%$) and $\psi'$ ($\sim 10\%$) decay. Since the latter melt earlier, this leads to the two-step suppression pattern [31, 32] already mentioned in section 2a: once we pass the deconfinement threshold, first the $J/\psi$'s originating from $\chi$ and $\psi'$ decay melt, and then, for sufficiently high energy density, also the directly produced $1S$ states start to disappear.

The statistics of the data shown in Fig. 15 are clearly not sufficient to test such details. However, the 1996 run of NA50 has provided more than five times as many $J/\psi$'s (about 275 000), and the new but still preliminary data look very interesting indeed [33]. In Fig. 18, we compare it to a form based on a first order deconfinement transition with critical bubble size formation [34]; we emphasize that this Figure is only meant to illustrate the present preliminary experimental pattern and the general theoretical behaviour due to a discontinuous onset of deconfinement. A more detailed theoretical study will have to wait for the final data analysis. We note here only that even for a discontinuous deconfinement onset, the $E_T - b$ smearing needed to relate theory to experiment will cause a considerable softening of the discontinuity, as seen in Fig. 18. Nevertheless, if such a step-like suppression pattern is confirmed by the final data, any hadronic comover explanation is excluded already on a qualitative level.
Figure 14: $J/\psi$ production in $A-B$ and $Pb-Pb$ collisions [26], compared to pre-resonance absorption in nuclear matter [23].

Figure 15: The $E_T$-dependence of $J/\psi$ production in $Pb-Pb$ collisions [26], compared to pre-resonance absorption in nuclear matter [23].
Figure 16: $J/\psi$ suppression by hadronic comover absorption compared to that by colour deconfinement.

Figure 17: Energy density profiles for central $A - B$ collisions above and below the deconfinement threshold.
Figure 18: Preliminary data on the $E_T$-dependence of $J/\psi$ production in $Pb-Pb$ collisions [33], compared to the pattern from a first order deconfinement transition [34].

5. Outlook and Summary

To establish that the observed anomalous $J/\psi$ suppression is indeed due to the onset of colour deconfinement, some further experimental studies are clearly necessary.

- First and foremost, it remains to be seen if the final analysis of the 1996 $Pb-Pb$ run corroborates the step-like onset of $J/\psi$ suppression. To determine the crucial variable for a deconfinement onset, one should then further study the effect for $A-A$ collisions using different collision energies and different $A$.

- Additional information will come from a study of the dependence of the anomalous suppression on the transverse momentum of the $J/\psi$. Initial state multiple scattering of the gluons fusing to form the $c\bar{c}$ pair leads to a broadening of the $P_T$ spectra from nuclear targets, compared to that in $p-p$ collisions. This ‘normal’ $P_T$ broadening accounts for all observed $J/\psi P_T$-distributions up to central $S-U$ collisions [35]. The melting process in anomalous $J/\psi$ suppression removes exactly those $J/\psi$’s (in the central tube of the interaction region) which came from gluons with most initial state scattering. This should lead to a characteristic turn-over of $P_T$ broadening [36], and it will be of great interest to see if such an anomalous $P_T$ behaviour of $J/\psi$ suppression is shown by the data.

- A further signal for a second step corresponding to the onset of dissociation of directly produced $J/\psi$’s would be provided by an increase in the ratio of $\psi'$ to $J/\psi$ production.
Finally, it is obviously challenging to check if an onset of colour deconfinement in the early stages of nuclear collisions, as are probed by charmonia, will have repercussions also on other observables.

Some of these questions may well be answered within the next few months, when the final analysis of the NA50 is presented. The search for colour deconfinement thus appears indeed to have reached a rather decisive stage.

In summary: statistical QCD predicts colour deconfinement in media of a sufficiently high density of constituents. Nuclear collisions are expected to lead to such conditions, and quarkonium states provide a probe to check the confinement/deconfinement status of the produced medium.

Up to central $S-U$ collisions, nucleus-nucleus data do not show any $J/\psi$ attenuation beyond ‘normal’ pre-resonance absorption in the nuclear matter of target and projectile. In contrast, $Pb-Pb$ collisions appear to lead with increasing centrality to an abrupt onset of a further ‘anomalous’ suppression. Such a sudden change in behaviour has no ‘conventional’ explanation; if confirmed in the final data analysis, it would indeed find a natural account only in terms of some form of critical behaviour.

Acknowledgements:

Much of what is presented here is based on joint work with D. Kharzeev and M. Nardi; I thank them for this fruitful and enjoyable collaboration. It is also a pleasure to thank M. Gonin, L. Kluberg and C. Lourenço for stimulating discussions of the NA38/NA50 data.

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