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Colour Confinement in Hadron-Nucleus Collisions

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Abstract:

We study the evolution of colour confinement in hadron-nucleus collisions and determine the resulting x_F dependence for the suppression of quarkonium production on nuclear targets. The Landau-Pomeranchuk effect applied to colour bremsstrahlung is shown to play a considerable role in the form of the suppression.

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The production of heavy quark ("quarkonium") resonances in hadron-nucleus collisions provides a crucial test of our understanding of colour confinement in QCD. The intrinsic masses of c and b quarks are sufficiently large to expect perturbative QCD to be applicable for the initial quark-antiquark formation process in most of phase space [1-3]. At the time of its formation, such a $Q\bar{Q}$ system is generally not colour neutral. Confinement certainly requires the final quarkonium state to be a colour singlet, but the time needed for the neutralisation of colour can be much shorter than the formation time for the physical resonance. Up to what time and size can a system remain coloured? The obvious way to measure the colour neutralisation scale is to send a $Q\bar{Q}$ system through extended matter susceptible to colour interactions and see when this matter becomes transparent to the probe. Since the passage of a nascent $Q\bar{Q}$ pair through a nucleus can bring it close to the coloured partonic constituents of several spatially separated nucleons, quarkonium production in $p-A$ collisions seems the ideal tool for such studies.

The argument may also be turned around. The suppression of quarkonium production can provide a signature for deconfined matter produced in high energy nucleus-nucleus collisions [4]. An essential prerequisite for this is an understanding of the colour confinement scale in the corresponding hadron-nucleus process. The matter in a nucleus is not deconfined, and to attest deconfinement, the life-time of the coloured state in nuclear collisions must therefore exceed that in $h-A$ collisions.

To discuss a specific case, let us consider J/ψ production, first in hadron-hadron collisions. We separate the process into a perturbative (gluon fusion) production stage (see Fig. 1a), resulting in a coloured $c\bar{c}$ pair. This is followed by the emission or absorption of a further gluon, leading to colour neutrality and the correct (3S_1) quantum numbers. The final step is the evolution of this nascent J/ψ into a fully grown physical charmonium resonance, with its given size and decay patterns. For J/ψ production at large transverse momenta, this process can be viewed as perturbative three-gluon fusion [5]; the third gluon can be attached to either projectile or target constituents. In the low P_T region, in which we are here primarily interested, the colour neutralisation process may be arbitrarily soft and can also involve the emission of several soft gluons; in any case, its effect cannot be calculated in a perturbative framework. We shall here assume a phenomenological extension into the non-perturbative soft-gluon regime, keeping the structure of the process as schematically illustrated in Fig. 1a.

In principle, the soft gluon could also be attached between the two hard ones, as shown in Fig. 1b. However, the soft gluon in this diagram couples to a quark which is far off-shell; this introduces an additional suppression factor in the amplitude, due to the extra quark propagator. In contrast, in the amplitude of Fig. 1a, the quark after the third gluon is kept essentially on mass shell by the J/ψ wave function; if the gluon is soft enough, the intermediate quark propagator is also close to its pole and hence no strong suppression of the amplitude occurs. The amplitude of Fig. 1a will therefore provide the main contribution in soft gluon emission.

The validity of gluon fusion dominance in quarkonium production is supported by two recent experimental observations [6]. The cross-sections for J/ψ production by π^+ , π^- , $\bar{p}$ and p beams of 300 GeV/c, incident on a Lithium target, are approximately

equal in most of phase space, indicating that gluon fusion indeed is the dominant mechanism. The measured ratio of ψ' to J/ψ production is correctly described by the ratio of the corresponding wave functions for on-shell $c\bar{c}$ systems [3]; since the geometric sizes of ψ' and J/ψ are quite different, this indicates that the interaction leading to colour neutralisation occurred well before the quarkonium states were fully formed.

Turning now to hadron-nucleus interactions, we note that the virtual colour-octet state in Fig. 1a will experience colour interactions in nuclear matter, provided it is present in the medium over a distance exceeding the characteristic scale of a single nucleon; the range in which these interactions will play a role depends on the characteristic lifetime of the coloured state and on the time between successive collisions in the medium.

To estimate how long the colour octet state will live, we first note that the J/ψ wave function keeps the charmed quarks close to their mass shells; we have thus $p^2 \simeq m_c^2$ (see Fig. 1 for the notation). The intermediate quark with four-momentum $k = p + q$ is off-shell by an amount

$$\Delta = (p + q)^2 - m_c^2 = 2pq, \quad (1)$$

where q is the momentum of the third gluon; in the spirit of the parton model, we take all gluons to be on-shell. The proper lifetime of the virtual state is given by the uncertainty principle as

$$\tau_0 \simeq |\Delta|^{-1/2}. \quad (2)$$

It is clear from eqs. (1,2) that the colour-octet state can propagate over a long distance only if the third gluon is soft enough. In the rest frame of the J/ψ , the lifetime of the colour-octet fluctuation,

$$\tau_0 \simeq (2m_c\omega)^{-1/2}, \quad (3)$$

is determined by the energy ω of the third gluon.

The properties of gluon propagators in the infrared region are still poorly known (see [7] for a recent discussion); it is nevertheless clear that gluons cannot propagate over large distances in a confining medium, be it the QCD vacuum or nuclear matter. Therefore one expects the existence of a cut-off factor ω_0 , with only the frequencies $\omega \geq \omega_0$ allowed. The cut-off parameter ω_0 should be at least of the order of $R_{QCD}^{-1} \simeq \Lambda_{QCD} \simeq 0.2$ GeV, which gives us

$$\tau_0 \simeq (2m_c\Lambda_{QCD})^{-1/2} \simeq 0.25 \text{ fm} \quad (4)$$

for the proper lifetime of the colour octet state. This estimate for the value of the colour octet $c\bar{c}$ lifetime is in good agreement with that extracted [8] from the data on J/ψ production in muon-nucleon interactions *

* The authors of [8] assume that confining force acting between colour octet states is the same as between quarks, and from the analysis of J/ψ rapidity distributions they then get the value $\tau_0 = 0.7 \pm 0.2$ fm. As we shall discuss below, the string tension for colour octet states is 9/4 times larger than that for colour triplets [9,10]; with this correction, the result of [8] becomes $\tau_0 \simeq 0.3 \pm 0.1$ fm, in reasonable agreement with our estimate (4).

Let us note at this point that the break-up of the $c\bar{c}$ pair into an open charm $D\bar{D}$ state,

$$\bar{\tau}_0 \simeq (2M_D - M_{J/\psi})^{-1} \simeq 0.31 \text{ fm} \quad (5)$$

is essentially of the same size. We can thus say more generally that after the formation of the coloured $c\bar{c}$ pair, a characteristic time of about 0.3 fm is needed before a colour-neutral object can appear.

In the rest frame of the target nucleus, in which the produced $c\bar{c}$ pair is generally fast, the proper time is Lorentz-boosted by the factor

$$\gamma = \frac{E}{M} = \left(1 + \frac{P^2}{M^2}\right)^{1/2}, \quad (6)$$

with

$$P = \frac{x_F s}{4m} \left(1 \pm \sqrt{1 + \frac{4M^2}{x_F^2 s}}\right), \quad (7)$$

and the positive (negative) sign for $x_F > 0$ ($x_F < 0$). In eqs. (6) and (7), M is the mass, E the lab energy and P the lab momentum of the quarkonium; s is the total incident energy in the hadron-nucleon center of mass system and m the nucleon mass. The Feynman variable is defined as $x_F = 2P_{cm}/\sqrt{s}$, where P_{cm} is the longitudinal momentum of the produced quarkonium in the projectile-nucleon center of mass.

The $c\bar{c}$ pair becomes colour-neutral after a time τ_0 , measured in its own rest frame. In the target rest frame, we have correspondingly the time $t_0 = \gamma\tau_0$; the distance (again measured in the target rest frame) which the pair travels during this time from its point of creation will be denoted by $s_0 = (P/M)\tau_0 = (E/M)t_0$.

We now come to the evolution of the $c\bar{c}$ system in the target. It depends crucially on the time between two successive interactions, compared to t_0 . The reason for this is the QCD-extension of the well-known Landau-Pomeranchuk effect [11]. An electron scattered in the field of an electric charge requires a certain characteristic time τ_0 for the resulting bremsstrahlung – a photon – to be formed and emitted. If during this time, the electron undergoes further interactions in the medium, the bremsstrahlung does not have time to develop and the clock measuring τ_0 is set back to zero after each interaction. As a consequence, an electron which in traversing the medium suffers n successive scatterings, at time steps less than τ_0 apart, will only radiate one photon, which is emitted when the electron leaves the medium. To see how this effect is relevant in our case, we picture the process shown in Fig. 1a in a slightly different way (Fig. 2a). The initial gluon g_1 fragments into a colour octet $c\bar{c}$ pair and subsequently interacts with the colour charge g_2 . This causes in general a deflection of the colour orientation of the $c\bar{c}$ state, leading as a consequence to the emission a third gluon (colour bremsstrahlung). We are interested here at the moment only in the case in which this second deflection results in a colour neutral state. Just as above in the electric bremsstrahlung, the emission of the colour bremsstrahlung g_3 requires a certain characteristic time τ_0 , and if the colour octet $c\bar{c}$ undergoes successive interactions, then these set the clock here also back to zero and the colour-neutralising gluon is emitted only when the $c\bar{c}$ leaves the medium. The occurrence of successive interactions in a dense medium thus retards the

emission of the bremsstrahlung and hence in our case the onset of colour neutralisation. As noted, the criterion for this effect is that the mean free path of the colour octet in the medium, which we shall denote by s_8 , is less than the distance s_0 the pair would travel before it can undergo colour neutralisation,

$$s_0 > s_8 = (n_0 \sigma_8)^{-1}. \quad (8)$$

Here $n_0 = 0.17 \text{ fm}^{-3}$ is standard nuclear density and σ_8 the cross section for interactions of the colour octet with the nuclear medium. What can be said about such interactions? The small size of the $Q\bar{Q}$ pair is certainly no longer any reason for weak interactions; even a pointlike colour charge will undergo strong interactions in nuclear matter. It is difficult to estimate of the size of σ_8 without referring to some particular model. To get some idea, let us assume that the overall colour neutrality is provided by a comoving light $q\bar{q}$ pair. The interaction cross section can then be defined for the colour neutral four quark $q\bar{q}Q\bar{Q}$ system. Its size is determined by the size of the light $q\bar{q}$ pair, so that the corresponding cross section should be the same as that for light $q\bar{q}$ mesons. We thus conclude that the $q\bar{q}Q\bar{Q}$ system should scatter on intranuclear nucleons with a cross section of the order of

$$\sigma_8 \simeq \sigma_{\pi N} \simeq 20 \text{ mb}, \quad (9)$$

which is about ten times larger than the cross section of a physical J/ψ on nucleons. This gives us in turn for the mean free path of a colour charge in nuclear matter the value $s_8 \simeq 3 \text{ fm}$.

The evolution of the $c\bar{c}$ thus leads to several distinct time or length scales; let us consider all as measured in the target rest frame. The $c\bar{c}$ system becomes colour neutral after a time t_0 , in which it has travelled a distance s_0 . It continues to expand, and after time $t_f \geq t_0$ (or distance s_f) it has become a resonance ($J/\psi, \psi'$) of final physical size. Competing with these scales is the time t_8 (or distance s_8) between two successive scatterings of the colour octet in the medium, and the time t_A (distance L_A) after which $c\bar{c}$ pair has left the target.

According to the relative values of these scales in different kinematic situations, we encounter several physically distinct regimes.

(1) The resonance regime.

When the distance travelled by the $c\bar{c}$ during the time needed for the formation of the physical quarkonium state is less than the intranuclear separation, the state will interact with its physical cross section over its entire path through the nucleus. This implies in particular that ψ' suppression must be stronger than J/ψ suppression, since the ψ' is larger than the J/ψ and hence has a larger absorption cross section. The resonance regime is thus fully attained when

$$s_f = \left(\frac{P}{E}\right) t_f = \left(\frac{P}{M}\right) \tau_f \geq z_0, \quad (10)$$

where z_0 is the average intranuclear distance. From the average nuclear density $\rho_0 = 0.17 \text{ fm}^{-3}$, we get $z_0 \simeq \rho_0^{-1/3} \simeq 1.8 \text{ fm}$, while nucleonic hard-core estimates give $z_0 \simeq 1.2 \text{ fm}$; we shall here use the latter value, since it is the smallest distance at

which the interaction with another nucleon can start. The J/ψ formation time can be determined through charmonium spectroscopy [12] as well as from the separation of J/ψ and ψ' masses [13], resulting in $\tau_f \simeq 0.4 - 0.5 \text{ fm}$; the $c\bar{c}$ pair thus needs about twice as long to become a fully established J/ψ as it needs to neutralise its colour. Using these values together with relations (6) and (7), we find that for energies $\sqrt{s} = 20 - 40 \text{ GeV}$, the physical resonance interaction regime is fully developed for $x_F \simeq -0.5$ or smaller.

(2) The colour regime.

When the distance travelled by the $c\bar{c}$ during the time needed to neutralise its colour exceeds the length L_A of its path through the medium, it will experience colour interaction throughout the entire target. If the initial gluon fusion occurs at a random point inside a nucleus of radius R_A , then $L_A = 3R_A/4$. For a standard heavy nucleus of $A = 200$ and with $R_A = 1.12A^{1/3}$, this gives us $L_A \simeq 5 \text{ fm}$. With a proper colour neutralisation time $\tau_0 = 0.25 \text{ fm}$, the relation

$$s_0 \geq L_A, \quad (11)$$

then implies that the $c\bar{c}$ will remain coloured on its entire path for $x_F \gtrsim 0.2$ at $\sqrt{s} = 20 \text{ GeV}$ and for $x_F \gtrsim 0$ at $\sqrt{s} = 40 \text{ GeV}$. In obtaining this estimate, however, we have ignored the Landau-Pomeranchuk effect, which makes the colour regime in fact start at much lower x_F . As soon as

$$s_0 \geq s_8, \quad (12)$$

the $c\bar{c}$ will undergo further and further scattering in the medium before it has had a chance to neutralise its colour, and the clock measuring the time until colour neutralisation is set back to zero. As the $c\bar{c}$ continues on its path, this effect will repeat itself again and again, and hence the $c\bar{c}$ will (on the average) traverse the entire medium as a colour octet. Using the value $s_8 \simeq 3 \text{ fm}$ obtained above together with eq. (12), we thus obtain that the $c\bar{c}$ will leave the target as colour octet for $x_F \gtrsim 0$ at $\sqrt{s} = 20 \text{ GeV}$ and for $x_F \gtrsim -0.2$ at $\sqrt{s} = 40 \text{ GeV}$.

The resulting resonance and colour interaction regimes are illustrated in Fig. 3. The two are separated by a transition region which for $\sqrt{s} \rightarrow \infty$ and with $\tau_f = 0.5 \text{ fm}$ becomes the interval $-0.55 \leq x_F \leq -0.27$; for $\sqrt{s} = 40 \text{ GeV}$, this is almost reached. In this transition regime, the system on its way through the nucleus is partially a small colour singlet, below the physical J/ψ size, and partially a colour octet. For $s_f > L_A$, which means for $x_F > 0$ at $\sqrt{s} = 20 \text{ GeV}$ and for $x_F > -0.1$ at $\sqrt{s} = 40 \text{ GeV}$, the $c\bar{c}$ will not have reached the full J/ψ size when it leaves the medium; hence for positive x_F , there should be no difference between J/ψ and ψ' nuclear suppression as effects in the transition region are colour transparency [14], gluon shadowing [15] and the onset of colour octet interactions [16]. We now want to consider in more detail the form of colour octet interactions and the resulting behaviour for $x_F \geq 0$, i.e., in the region in which the colour neutralisation distance exceeds L_A .

One crucial feature of the interaction of the colour octet $c\bar{c}$ in the target medium is that it cannot lead to a break-up of the pair. We had seen above that the break-up

process with subsequent $D\bar{D}$ production takes the same length of time ($\bar{\tau}_0 \simeq 0.3$ fm) as colour neutralisation; it is also of the same type (colour bremsstrahlung), as seen in Fig. 2b. Successive interactions of the $c\bar{c}$ in the medium thus leave the pair together for the same reason they leave it coloured, and it can fragment into a $D\bar{D}$ only when it leaves the medium. Thus an increase in the number of interactions will not increase the suppression of the J/ψ .

The second decisive feature is that the coloured pair in traversing the nucleus runs into bubbles of physical vacuum situated between the individual nucleons. These collisions "stretch a string" between the pair and the nearest nucleon, and the stretching of this string over the mean free path $s_8 \simeq 3$ fm of the colour octet inside the target causes the $c\bar{c}$ to lose energy step by step along its path. If the pair initially has the momentum P , then in the collision this will be shifted by

$$P \rightarrow (P - \kappa s_8), \quad (13)$$

where κ denotes the hadronic string tension. The next collision will have the same consequence, and on leaving the target after n successive collisions, the momentum of the pair has been degraded by $n\kappa s_8 = \kappa(L_A - z_0)$. This means that a J/ψ observed at a given x_F comes from a pair originally produced at the higher value x_F/α , with

$$\alpha \equiv \left(\frac{P - \kappa(L_A - z_0)}{P} \right). \quad (14)$$

The resulting normalised x_F -distribution $F_A(x_F) \equiv (1/A)(d\sigma_{hA \rightarrow J/\psi}/dx_F)$ for J/ψ production in $h - A$ collisions thus becomes

$$F_A(x_F) = W_0 F_0(x_F) + (1 - W_0) \left(\frac{F_0(x_F/\alpha)}{\alpha} \right) \Theta(1 - x_F/\alpha). \quad (15)$$

Here $F_0(x_F)$ is the x_F -distribution in $h - p$ collisions and

$$W_0 \simeq \exp\{-(L_A - z_0)/s_8\} \quad (16)$$

gives the probability that the $c\bar{c}$ pair will not undergo any scattering on its way out of the target; thus $W_0 = 1$ implies no nuclear J/ψ suppression. The second term in eq. (15) gives the effect of the n coherent scatterings in the medium, with the resulting shift in x_F . The factor $1/\alpha$ assures that the distribution remains normalised. Note that the fate of the $c\bar{c}$ pair for $x_F \geq 0$ as described here does not imply any integrated J/ψ suppression: the production is simply shifted from larger to smaller x_F .

To determine α , we need the hadronic string tension, which in lattice gauge theory is found to be around 1 GeV/fm for the force between static quarks, i.e., between colour triplets. The confining force between colour octet objects has been studied in lattice QCD [9] as well as in analytical approaches[10]. In the latter case the string tension appears to be proportional to the quadratic Casimir operator of $SU(N_c)$ colour group; for the ratio of the "adjoint" to "fundamental" string tensions we have

$$\frac{\kappa_{\text{adj}}}{\kappa_{\text{fund}}} = \frac{2N_c^2}{N_c^2 - 1}. \quad (17)$$

In the large N_c limit, the ratio (17) becomes 2, in agreement with $1/N_c$ factorization [17]; for real QCD, $N_c = 3$ and hence the string tension between colour octet states is 9/4 times larger than that between quarks. This property of the confining force was in fact observed in lattice Monte Carlo calculations [9]. We use therefore in the following the value $\kappa \equiv \kappa_8 = (9/4)\kappa_3 \simeq 9/4$ GeV/fm.

Dividing eq. (15) by the x_F -distribution for $h-p$ collisions, $F_0(x_F)$, gives us the experimentally measured ratio of $h - A$ to $h - p$ distributions,

$$R_{A/p}(x_F) = W_0 + (1 - W_0) \left\{ \frac{F_0(x_F/\alpha)}{\alpha F_0(x_F)} \right\} \Theta(1 - x_F/\alpha). \quad (18)$$

We use the values of the different parameters as obtained above, i.e., $L_A \simeq 5$ fm for a typical heavy nucleus of $A = 200$, $s_8 \simeq 3$ fm for a colour octet interaction cross section of 20 mb, an intranuclear distance $z_0 \simeq 1.2$ fm, and an octet string tension $\kappa \simeq 9/4$ GeV/fm. This gives us the dependence on x_F shown in Fig. 4 for three different beam energies. We note that on nuclei, compared to nucleons as target, J/ψ production is shifted to lower x_F , essentially because of the momentum degradation of the colour octet in passing through the medium. This momentum loss is $\Delta P = \kappa(L_A - z_0) \simeq 8.6$ GeV and thus fixed (see [18] for a similar discussion in case of a coloured parton passing through a nuclear medium). As a consequence, the relative shift in x_F decreases with increasing beam energy, and hence the effect of the medium disappears for $\sqrt{s} \rightarrow \infty$, when $R_{A/p} \simeq 1$ for $0 < x_F < 1$. While the general pattern seen in Fig. 4 agrees with that found experimentally [19,20], the colour interactions do not lead to the observed suppression. As has been noted [15], the actual suppression well outside the resonance region is largely due to gluon shadowing. In this work, such shadowing was found to have the form

$$R_{A/p}^h = 1 + a \ln x_2 \ln A, \quad (19)$$

where

$$x_2 = \frac{1}{2} \left(\sqrt{4M^2/s} + x_F^2 - x_F \right) \quad (20)$$

denotes the Bjorken x of the target. Assuming all x_F -dependence in the region $x_F \leq 0.5$ to be due to gluon shadowing gave $a \simeq 0.021$. We saw in the present study that colour octet interactions do give a further x_F -dependence in the region in question. This will somewhat weaken the shadowing and thus bring it closer to the behaviour found for quark shadowing [21]. We shall therefore assume the same slopes in $(\ln x_2 \ln A)$ for gluon shadowing as found for quark shadowing in the relevant x_2 region; this suggests $a = 0.018$ instead of 0.021. In Fig's. 5 and 6 we compare the resulting suppression

$$S_{A/p}(x_F) \equiv R_{A/p} \times R_{A/p}^h, \quad (21)$$

as function of x_F to data from pion and proton beams at different energies incident on heavy targets. The agreement is seen to be quite good, in spite of the rather rough way we have treated the nuclear geometry of the problem. At the lowest pion beam energy ($P = 150$ GeV, i.e., $\sqrt{s} = 17$ GeV), we note that the predictions appear to somewhat overshoot the data. This is not unexpected, since outside the resonance regime, there

will still be some absorption of the nascent $c\bar{c}$ pairs. This effect can take place only as long as

$$\left(\frac{P}{M}\right) \bar{\tau}_0 \leq L_A, \quad (22)$$

where $\bar{\tau}_0$ is the break-up time defined in eq. (5). For $\sqrt{s} = 17$ GeV, this region extends to $x_F \simeq 0.25$; for higher energies, it moves to lower x_F . Such absorption, as well as the proper nuclear geometry, must be taken into account in a really quantitative study of the problem; this is underway.

We conclude here that the combination of colour octet interactions in a nuclear medium, taking into account the Landau-Pomeranchuk effect, and nuclear gluon shadowing of the same functional form as found for quarks, accounts quite well for the observed J/ψ suppression in $p - A$ and $\pi - A$ interactions.

Let us note briefly that the picture proposed here is consistent with colour transparency [14]. According to QCD factorisation, the cross-section for the hadroproduction of heavy quarks is given as a convolution of parton structure functions and the parton fusion cross section. If the momentum transfer in the parton fusion process is large enough (hadroproduction at high transverse momentum), one expects that a colour-neutral quark pair is produced in an almost pointlike state and hence final state nuclear effects should be small (colour transparency). On the other hand, if heavy quark pairs are produced in colour octet states, they interact immediately, so that no colour transparency is expected. Nevertheless, there is no contradiction here, since the lifetime τ_0 of the colour octet state is inversely proportional to the mass of the heavy quark m_Q (see Eq.(4)). Therefore, when $m_Q \rightarrow \infty$ (which implies an increase of the momentum transfer in parton fusion), $\tau_0 \rightarrow 0$ and the heavy quark pair traverses the nucleus in a colour neutral, pointlike state: colour transparency is recovered. For finite $m_Q = m_c(m_b)$ and/or large x_F , however, the pair traverses the nucleus in a coloured state with string interaction, and hence colour transparency does not set in. These considerations also throw some light on the relation between perturbative QCD and the string picture (see [22] and further references given there) of hadroproduction on nuclei.

In closing, we want to comment briefly on the problem of $Q\bar{Q}$ production in the high energy nucleus-nucleus collisions. If such collisions lead to the production of a deconfined medium, how this would influence the production of quarkonia according to the picture presented here? To answer this question we note that in deconfined matter, in contrast to the nuclear matter considered so far, gluons can propagate over large distances, with $R_g^{\text{dec}} \sim 5 - 10$ fm, as determined by the size of the system. The cut-off on the energy of the gluon can therefore be as low as $\omega_0 \simeq 20$ MeV. This value leads with eq. (3) in the case of $c\bar{c}$ states to

$$\tau_0^{\text{dec}} \simeq 0.8 \text{ fm} \quad (23)$$

for the proper lifetime of the colour octet state. While in nuclear matter the time needed to neutralise the colour of the $c\bar{c}$ state was about one-half that needed to form the physical resonance, it has now become about $2\tau_f$. This in turn implies that at the time when colour neutralisation would normally set in, the c and the $\bar{c}$ have separated

by twice the distance they are apart in the physical resonance. Hence in general they can no longer bind to form a J/ψ , and we thus confirm on a microscopic basis the earlier claim [4] that the suppression of quarkonium production can provide a signature for deconfinement. It thus appears that charmonium production in $h - A$ collisions (and for $x_F \geq 0$) tests the fate of a colour octet $c\bar{c}$ in a confining medium (the target nucleus), while such production in high energy nucleus-nucleus collisions (and $x_F \simeq 0$) probes its fate in a deconfining medium.

References:

- 1) H. Fritzsch, Phys. Lett. B67 (1977) 217.
- 2) M. Glück and E. Reya, Phys. Lett. B79 (1978) 453;
E. L. Berger and D. Jones, Phys. Rev. D23 (1981) 1521.
- 3) R. Baier and R. Rückl, Z. Phys. C19 (1983) 251.
- 4) T. Matsui and H. Satz, Phys. Lett. B178 (1986) 416.
- 5) L. Clavelli, P. H. Cox, B. Harms and S. Jones, Phys. Rev. D32 (1985) 612.
- 6) L. Antoniazzi et al., Phys. Rev. D46 (1992) 4828; L. Antoniazzi et al., Phys. Rev. Lett. 70 (1993) 383.
- 7) D. Zwanziger, Nucl. Phys. B378 (1992) 525.
J. R. Cudell and D. A. Ross, Nucl. Phys. B359 (1991) 247.
C. Bernard, A. Soni and K. Yee, Nucl. Phys. B20 (Proc.Suppl.) (1991) 410.
- 8) M. Erdmann and T. Dreyer, Phys. Lett. B260 (1991) 232.
W. Busza, T. Dreyer and M. Erdmann, Z. Phys. C48 (1990) 167.
- 9) J. Ambjorn, P. Olesen and V. Peterson, Nucl. Phys. B240 (1984) 189;
N. A. Campbell, I. H. Jorysz and C. Michael, Phys. Lett. B167 (1986) 91;
I. J. Ford, R. H. Dalitz and J. Hoek, Phys. Lett. B208 (1988) 286.
- 10) H. G. Dosch and Yu. A. Simonov, Phys. Lett. B205 (1988) 339;
Yu. A. Simonov, Nucl. Phys. B307 (1988) 512.
- 11) L. D. Landau and I. Ya. Pomeranchuk, Dokl. Akad. Nauk SSSR 92 (1953) 535, 735;
M. L. Ter-Mikaelyan, Zh. Eksp. Teor. Fiz. 25 (1953) 289, 296;
A. B. Migdal, Zh. Eksp. Teor. Fiz. 32 (1957) 633 (Sov. Phys. JETP 5 (1957) 527);
E. L. Feinberg, Zh. Eksp. Teor. Fiz. 50 (1966) 202 (Sov. Phys. JETP 23 (1966) 132).
- 12) F. Karsch and H. Satz, Z. Phys. C51 (1991) 209.
- 13) D. E. Kharzeev, Inst. Phys. Conf. Series 124 (1992) 229; Nucl. Phys. A, in press.
- 14) S. J. Brodsky and A. H. Mueller, Phys. Lett. B206 (1988) 685.
- 15) S. Gupta and H. Satz, Z. Phys. C55 (1992) 391.
- 16) G. Piller, J. Mützbauer and W. Weise, Z. Phys. A343 (1992) 247.
- 17) J. Greensite and M. B. Halpern, Phys. Rev. D27 (1983) 2545.
- 18) S. J. Brodsky and P. Hoyer, Phys. Lett. B298 (1993) 165.
- 19) J. Badier et al., Z. Phys. C20 (1983) 101.
- 20) D. M. Alde et al., Phys. Rev. Lett. 66 (1991) 133;
D. M. Alde et al., Phys. Rev. Lett. 66 (1991) 2285.
- 21) J. J. Aubert et al., Phys. Lett. B123 (1983) 275;
A. C. Benvenuti et al., Phys. Lett. B189 (1987) 483;
P. Amaudruz et al., Z. Phys. C51 (1991) 387.
- 22) M. Gyulassy and M. Plümer, Nucl. Phys. B346 (1990) 1;
B. Z. Kopeliovich, Phys. Lett. B243 (1990) 141;
K. Borekov, A. Capella, A. Kaidalov, J. Tran Thanh Van, Phys. Rev. D47 (1993) 919.

Figure Captions

Fig. 1 J/ψ production through gluon fusion.

Fig. 2 Colour bremsstrahlung leading to J/ψ production (a) and to a $D\bar{D}$ pair (b).

Fig. 3 Different interaction regimes for J/ψ production in $p - A$ collisions at two different energies: resonance regime (RR) and colour regime (CR) with (LP) and without (NO LP) Landau-Pomeranchuk effect.

Fig. 4 The effect of colour octet interactions on the x_F distribution for J/ψ production in $p - A$ collisions at three different energies.

Fig. 5 The x_F distributions for J/ψ production in $\pi^- - A$ interaction at three different energies [23], compared to the prediction from colour interactions (solid line).

Fig. 6 The x_F distributions for J/ψ production in $p - A$ interaction at two different energies [20], compared to the prediction from colour interactions (solid line).

Figure 1

Figure 2

Figure 3

Figure 4

Figure 5

Figure 6