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DECONFINEMENT: CONCEPT, THEORY, TEST*

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Abstract

We survey the physical ideas leading to deconfinement, discuss the transition from hadron matter to quark-gluon plasma in statistical QCD, and give a confinement/deconfinement probe for strongly interacting matter.

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INTRODUCTION

The study of strongly interacting matter is a fascinating subject for a variety of reasons. It is needed to describe the first few microseconds in the evolution of our universe, and it provides the equation of state for stellar matter at high density. It gives us the statistical closure of whatever strong interaction dynamics we consider, from pion exchange to QCD. As such, it was and remains a laboratory to investigate the critical behaviour arising from new forms of dynamics. Hagedorn's ultimate temperature of hadronic matter [1] is one particularly striking result obtained in this laboratory – and one which did much to further our understanding of dense matter. In the past decade, two new developments made strong interaction thermodynamics even more attractive.

On the theoretical side, large-scale computer simulations allow for the first time in physics an *ab initio* calculation of the thermodynamic behaviour of an actual physical system. From QCD as input dynamics, we can today evaluate directly, i.e., without intermediary models, the states of matter and the critical behaviour of strongly interacting systems. The most important result which has so far come out of these studies is a quantitative confirmation of the idea that matter at sufficiently high density will undergo a transition to a state of deconfined quarks and gluons.

On the experimental side, an extensive program of high energy nuclear collision studies was initiated, with the specific aim of investigating the critical behaviour and the new deconfined state of strongly interacting matter under controllable laboratory conditions. The collision of two heavy nuclei will, as we hope, provide droplets of such matter sufficiently large, hot and thermalised to form the quark-gluon plasma predicted by statistical QCD. The experiments carried out so far, using beams of rather light nuclei on heavy targets, have shown that the analysis of very high multiplicity final states is feasible, and they have led to a number of results which seem to indicate that we are indeed on the road to strong interaction thermodynamics.

In this survey of deconfinement, we will first have a look at the physical concepts which led to the phenomenon, and then turn to statistical QCD and to the results obtained from it by computer simulation on the lattice. Finally we will indicate how one can test if the matter produced in high energy nucleus-nucleus collisions indeed consists of deconfined quarks and gluons.

COLOUR SCREENING IN DENSE MATTER

An increase in the energy density of a pion gas leads not only to the production of more pions, but also to abundant resonance production. One way to accommodate this is to consider the thermodynamics of an ideal resonance gas [2]. Its partition function is

$$\ln Z_R(T, V) = \frac{VT^3}{\pi^2} \int dm \rho(m) e^{-m/T}, \quad (1)$$

where $\rho(m)$ describes the resonance mass spectrum; for $\rho(m) = \delta(m - m_\pi)$, we get back the ideal pion gas. To fix it in general, we need an input from hadron dynamics. Hagedorn [1] used a composition law, the statistical bootstrap model, requiring resonances to be composed of resonances in a universal way. It yields

$$\rho(m) \sim m^a e^{bm}, \quad (2)$$

with the constant a dependent on the details of the model, while b is fixed by the lowest hadron mass, m_π , giving $b \simeq 150$ MeV. The most detailed model for hadronic interactions, the dual resonance model [3], assumes the interaction amplitude to be built up by resonances in all channels; it leads to the same form (2). In this case, b is determined by the Regge slope governing the spin degeneracy of resonances; the resulting value agrees remarkably well with that from the statistical bootstrap [4]. A similar form had previously also been obtained through geometric arguments [5]; here b was determined through the size of the basic hadrons, leading again to a value near 150 MeV.

Inserting the spectral form (2) into the partition function (1), we note that the integral diverges for $T > T_H \equiv 1/b$. What is the significance of this temperature? Hagedorn had initially proposed it to be the ultimate temperature of matter, since for the values of a he had obtained, $T \rightarrow T_H$ was possible only for an energy density $\epsilon \rightarrow \infty$. However, it was noted soon afterwards [6] that for slightly smaller values of a , ϵ remained finite at $T = T_H$; only higher derivatives of the partition function diverged, suggesting a transition to a new phase of matter.

Whatever interpretation we choose, we are left with the conclusion that there is an intrinsic limit to hadron thermodynamics; apparently something must happen at the Hagedorn temperature $T_H \simeq 150$ MeV. Today we consider hadrons as bound states of quarks, and it therefore seems natural to associate T_H to a transition from hadronic to quark matter.

This raises some questions. What happens to colour confinement in such a transition? How can the binding of quarks be dissolved, when the confinement potential $V_{\text{conf}}(r) \sim r$ diverges as $r \rightarrow \infty$?

The mechanism responsible for colour deconfinement is charge screening in a dense medium. The Coulomb potential between two electric charges becomes Debye-screened in a dense ionised gas,

$$\frac{e^2}{r} \rightarrow \left(\frac{e^2}{r} \right) e^{-\mu_D r}, \quad (3)$$

with a screening radius $r_D = 1/\mu_D$. For the same reason, the colour potential will experience screening in a medium of colour charges,

$$\sigma r \rightarrow \left(\frac{1 - e^{-\mu r}}{\mu r} \right) \sigma r, \quad (4)$$

with $r = 1/\mu$ for the colour screening radius. The effect of this screening is seen in Fig. 1. With increasing colour charge density, μ increases (just as μ_D does for an increasing density of electric charges), and so eventually any bound state is dissolved. At low density, we thus have a *colour insulator*: hadronic matter, consisting of colour neutral bound quark states. At high density, this turns into a *colour conductor*: the quark-gluon plasma, whose constituents are unbound coloured quarks and gluons.

Figure 1. Effect of screening on a confining potential.

Colour deconfinement thus seems to be the QCD analogue of the insulator-conductor transition in QED. And if we assume the transition to occur when the screening radius has reached the typical hadronic scale of 1 fm, we find [7] a transition temperature of about 150 MeV.

What happens to the mass of the constituents in this transition? We know that the effective mass of the electron in a metal is not its vacuum mass – it becomes modified by medium effects (lattice vibrations, electron cloud). For strong interactions, with “bare” quark masses $m_q \simeq 0$ in the QCD Lagrangian, we have in the confined phase constituent quarks with an “effective” mass $m_q^{\text{eff}} \simeq m_p/3 \simeq m_\rho/2$. In this confined state, the chiral symmetry of the Lagrangian with massless fermions must therefore be spontaneously broken. At sufficiently high density, we expect the gluon dressing to come off, turning massive constituent quarks into massless current quarks and thereby restoring chiral symmetry.

For a Lagrangian with $m_q=0$, we thus expect a shift in the effective quark mass leading to chiral symmetry restoration, as a further phase transition in strongly interacting matter. Deconfinement and chiral symmetry restoration can, but need not coincide. In atomic physics, we can have an insulator turning with increasing density into a superconductor, which for still higher density becomes a conductor. The first transition, decoupling electron-ion binding, corresponds to deconfinement; however, the liberated electrons don’t remain free, but couple to Cooper pairs. In the second

transition, the Cooper pairs are broken up and the electrons obtain their “metallic” effective mass. In QCD, we could correspondingly have hadronic matter, constituent quark matter, and the quark-gluon plasma. Since constituent quarks, just like Cooper pairs, are probably easy to break up by thermal motion, such a constituent quark state may well exist only at low temperatures and high baryon number density. At zero baryon number density, deconfinement and chiral symmetry restoration occur at the same temperature, as lattice calculations (see the next section) show.

The schematic phase diagram of strongly interacting matter, as it emerges from these considerations, is shown in Fig. 2.

Figure 2. Schematic phase diagram for strongly interacting matter.

STATISTICAL QCD AT FINITE TEMPERATURE

Statistical QCD allows a *direct* determination of the phase structure of strongly interacting matter, starting from the QCD Lagrangian as dynamical input,

$$\mathcal{L} = -\frac{1}{4}(\partial_\mu A_\nu^a - \partial_\nu A_\mu^a - gf_{bc}^a A_\mu^b A_\nu^c)^2 - \sum_f \bar{\psi}_\alpha^f (i\gamma^\mu \partial_\mu + m_f - g\gamma^\mu A_\mu)_{\alpha\beta} \psi_\beta^f, \quad (5)$$

with m_f for the masses of N_f current quark species. The basis of thermodynamics is the QCD partition function, most conveniently written in Euclidean form as functional integral over gluon (A) and quark (ψ) fields,

$$Z(T, V) = \int DA D\psi D\bar{\psi} e^{-S(A, \psi, \bar{\psi})}, \quad (6)$$

where the action

$$S(A, \psi, \bar{\psi}) = \int_0^{1/T} d\tau \int_V d^3x \mathcal{L}(\tau = ix_0, \vec{x}) \quad (7)$$

is calculated directly from the Lagrangian $\mathcal{L}$. Derivatives of $Z(T, V)$ give us the thermodynamic quantities of interest, such as the energy density ϵ ,

$$\epsilon(T) = \left(\frac{T^2}{V} \right) \left(\frac{\partial \ln Z}{\partial T} \right)_V \quad (8)$$

or the pressure P ,

$$P(T) = T \left(\frac{\partial \ln Z}{\partial V} \right)_T. \quad (9)$$

Equations (5) - (7) thus constitute the complete basis of statistical QCD. We are left “only” with the question of how to carry out the calculation of the partition function (6/7).

So far, the computer simulation of the lattice formulation [8] is the only generally viable method for this, applicable in particular also in the critical region, where perturbative methods break down. Lattice QCD is obtained through three steps. First, we replace the space-time continuum $\{\tau = ix_0, \vec{x}\}$ in eq. (7) by a lattice of $N_\tau \times N_\sigma$ discrete points, with $\{\tau = N_\tau a; x_i = N_\sigma a, i = 1, 2, 3\}$ and an isotropic lattice spacing a ; hence $T = 1/(N_\tau a)$, $V = (N_\sigma a)^3$. Next, the quark and antiquark field integrations are carried out, which is possible since they enter the integral in the form $\exp(-\bar{\psi} Q_F \psi)$; this leads to a factor $(\det Q_F)$ in the remaining integral over the gluon fields A [9]. In the third step, these integrations are changed to the (compact) variables $U = \exp(ixA)$, where U denotes a matrix of the SU(3) colour gauge group of QCD; such variables U are associated to the links between any two points on the lattice. After these steps, the partition function (6) has the structure

$$Z(T, V) = \int \prod_{\text{links}} dU e^{-S(U)} \quad (10)$$

of a “generalized spin problem”, something like an Ising model partition function, with SU(3) matrices U instead of spins, and with these on lattice links instead of lattice sites. The QCD action $S(U) = S_G(U) + S_F(U)$ consists of the gluon part $S_G(U) \sim (1 - UUUU)$ and the quark part $S_F(U) \sim \log(\det Q_F)$. The product of four “spins” U on a closed loop of four links (“plaquette”) in the gluon action assures local gauge invariance [10], which the Ising action with its product of two adjacent spins does not have. This completes the lattice formulation. The results which we want to obtain from it are supposed to hold in the continuum limit, in which the number of lattice points becomes infinite while the lattice spacing simultaneously goes to zero, keeping T and V constant.

Now we want to study the critical behaviour arising in statistical QCD. Since phase transitions are in statistical physics quite generally associated with changes in symmetry, we look for the symmetries of the QCD Lagrangian $\mathcal{L}$ or of the QCD action $S(U)$ in eq. (10).

Consider the limit of very heavy quarks: for $m_f \rightarrow \infty$, $S_F \rightarrow 0$, and we get the thermodynamics of pure SU(3) gauge theory. This pure gluon system already contains many of the crucial features of QCD, quite in contrast to QED, where the pure photon system forms an ideal boson gas. The reason for the difference is the term $gA_\mu A_\nu$ in eq. (6), which allows gluons, unlike photons, to interact directly, without intermediate fermions.

If we now carry out the “spin flip”

$$U(n_0, \vec{n}) \rightarrow U'(n_0, \vec{n}) = zU(n_0, \vec{n}) \quad \forall \vec{n}, \quad (11)$$

where $z = \exp(i\pi r/3)$, $r = 0, 1, 2$, is an element of the center Z_3 of the SU(3) gauge group, then this leaves the action $S_G(U)$ invariant. Such an operation is just the analog of flipping all spins $s_i \rightarrow -s_i$ in the Ising model, where it also leaves the Ising

action invariant. However, the “generalized spin”

$$L(U) \equiv \prod_{n_0=1}^{N_\tau} U(n_0, \vec{n}) \quad (12)$$

is “flipped” by this operation: $L(U') = zL(U)$. Its average, $\langle L \rangle$, therefore measures a possible spontaneous breaking of the Z_3 symmetry of the state of matter in which the system is, just as the average spin $\langle s \rangle$ in the Ising model indicates if the Ising system is in an ordered state with spontaneous magnetization or in a disordered state with $\langle s \rangle = 0$.

The transition from unbroken to broken Z_3 symmetry corresponds in fact to deconfinement. The order parameter $\langle L \rangle$ can be related to the potential $V(r)$ between a static quark and antiquark in the limit of infinite separation,

$$\langle L \rangle \sim \lim_{r \rightarrow \infty} e^{-V(r)/T}. \quad (13)$$

For a confined state, $V(r)$ diverges for $r \rightarrow \infty$, so that $\langle L \rangle = 0$; therefore Z_N -symmetry signals confinement. In a deconfined state, $V(r)$ remains finite as $r \rightarrow \infty$, since colour screening cuts out the diverging (and hence confining) long distance part of the potential; thus the spontaneous breaking of the Z_N symmetry indicates deconfinement. Therefore the “generalized spin” $\langle L \rangle$ (in lattice gauge theory usually referred to as Polyakov loop) constitutes the theoretical probe for confinement or deconfinement. As noted, this is strictly true only in the limit of infinite quark mass, i.e., in pure $SU(N)$ gauge theory. For finite quark mass, thermal pair production leads to string breaking and hence to a finite $\langle L \rangle$ for all $T \neq 0$. This is quite similar to the actual insulator-conductor transition, where thermal ionisation also prevents the conductivity from vanishing in the insulator phase for $T \neq 0$. In both cases, however, the order parameter remains exponentially small in the “symmetric” phase: $\langle L \rangle \sim \exp(-m_h/T)$, with $m_h \sim m_\rho$ as typical hadron mass, corresponds to the conductivity $\sigma \sim \exp(-\Delta E/T)$, where ΔE denotes the ionisation energy.

We now turn to the already mentioned shift in the mass of the constituents in connection with deconfinement and consider the Lagrangian in the limit of vanishing quark mass, $m_f = 0$. The four-spinor fields for the quarks then decompose into a direct sum of a left-handed and a right-handed two-spinor, $\psi^{(4)} \rightarrow \psi_L^{(2)} + \psi_R^{(2)}$. The chiral “rotation” between these,

$$\psi \rightarrow \psi' = e^{i\alpha\gamma_5}\psi, \quad (14)$$

leaves invariant the $m_f = 0$ Lagrangian, since $\bar{\psi}'\gamma_\mu\psi' = \bar{\psi}\gamma_\mu\psi$. However, the mass term in the original $m_f \neq 0$ Lagrangian would not remain invariant, since $\bar{\psi}'\psi' = \bar{\psi}e^{2i\alpha\gamma_5}\psi \neq \bar{\psi}\psi$. Hence $\langle \bar{\psi}\psi \rangle$ measures if the system is in a chirally symmetric state or not. If $\langle \bar{\psi}\psi \rangle \neq 0$, the chiral symmetry of the state is spontaneously broken, the quarks have acquired a non-vanishing effective mass. For $\langle \bar{\psi}\psi \rangle = 0$, the chiral symmetry is restored and the constituents of the systems are quarks with a vanishing effective mass, i.e., the current quarks of the massless Lagrangian. In the real world, the quarks in the Lagrangian cannot be massless, since the pion mass $m_\pi \sim m_q^2$ is not zero, so that we have approximate chiral symmetry only.

Statistical QCD thus leads to two critical phenomena, deconfinement and chiral symmetry restoration. Deconfinement is associated to a global Z_3 symmetry of the Lagrangian for $m_f \rightarrow \infty$, with $\langle L \rangle$ as order parameter; the chiral symmetry of the

Lagrangian for $m_f \rightarrow 0$ is measured for a given state of the system by the order parameter $\langle \bar{\psi}\psi \rangle$. Although in the real world both these symmetries are only approximate, we nevertheless believe that their remnant effects will show up in the transition from low to high density behaviour of strongly interacting matter.

So the basic quantities to be studied in statistical QCD are the deconfinement measure $\langle L \rangle$, the chiral symmetry measure $\langle \bar{\psi}\psi \rangle$, the energy density ϵ , the pressure P , and thermodynamic quantities derived from these, such as entropy, specific heat, and others. The actual calculation of these “observables” is highly non-trivial, since we are studying a relativistic field theory for an interacting system near a critical point, where perturbative methods are not applicable. As already mentioned, the only viable evaluation method in this region is the computer simulation [11] of the lattice formulation of the problem [8]: one creates on a sufficiently large and fast computer a “world according to QCD” and then “measures” in this world the quantities of interest; for a recent review, see [12]. One draw-back of this method is that it can be applied up to now only to systems of zero over-all baryon number density, i.e., only to hot matter at $n_B = 0$, not to dense matter at $T = 0$; another is that its precision is limited by computer size and speed. Nevertheless, it provides us with the unique chance to calculate the critical behaviour directly from the underlying fundamental theory, and the mentioned shortcomings will hopefully be removed or reduced in the future.

The main results from the computer simulation of lattice QCD at finite temperature are summarized in Figs. 3 - 4, which were obtained in recent calculations for two light (staggered) quark species on $8^3 \times 4$ and $16^3 \times 4$ lattices [13]. In fig. 3, we see abrupt changes in the deconfinement measure $\langle L \rangle$ and the chiral measure $\langle \bar{\psi}\psi \rangle$ at the same temperature $T = T_c \simeq 150$ MeV; hence for $n_B = 0$, the two critical phenomena coincide. We also note that neither order parameter is ever really zero; for $\langle L \rangle$ this is due to finite lattice size and finite quark mass, for $\langle \bar{\psi}\psi \rangle$ due to the non-zero quark mass. In spite of this, both quantities show clear transition signals and thus allow the definition of a critical temperature T_c . In Fig. 4, we see that at this temperature, the dimensionless energy density measure ϵ/T^4 increases abruptly from a value near that of an ideal pion gas ($\simeq 1$) to one near that of an ideal quark-gluon plasma ($\simeq 20$). This increase in the energy density $\Delta\epsilon$ represents something like the “latent heat of deconfinement”. We also observe that the pressure grows with temperature at a much slower rate, and the ideal gas relation $\epsilon = 3P$ appears to become fulfilled only for $T \geq T_c$. Hence in the region $T_c \leq T \leq 2T_c$, there are definite interaction effects in the plasma. In spite of a number of model considerations (see e.g. [14] and further references there), their origin still constitutes at present one of the most interesting open problems in finite temperature lattice QCD.

We should note here that the value of the critical temperature in physical units ($T_c \simeq 150$ MeV) is obtained by calculating T_c in units of the ρ -mass and then using the experimental value of this mass.

Such a value of T_c means that an energy density $\epsilon_c \simeq 1.0$ GeV/fm³ is needed in order to reach deconfinement. Energy densities of this magnitude existed in the early universe for the first 10 - 20 microseconds after the big bang, and it is hoped to obtain them in small volumes and for short times in high energy nuclear collisions. Abundant multiple scattering among the incident nucleons and the many secondaries produced in the collisions is expected to provide a rapid increase in entropy sufficient to reach thermal equilibrium. And for high enough incident energies (15 - 20 GeV cms energy per nucleon-nucleon collision), the initial energy density will be above the deconfinement threshold of 1 GeV/fm³.

Figure 3. Temperature dependence of deconfinement measure $\langle \text{Re } L \rangle$ and chiral symmetry measure $\langle \bar{\psi}\psi \rangle$, from [13].

Figure 4. Temperature dependence of energy density ϵ and pressure P , from [13].

CONFINEMENT/DECONFINEMENT IN NUCLEAR COLLISIONS

The ultimate aim of high energy nuclear collisions is thus to produce and study the predicted deconfined state of strongly interacting matter. In whatever state matter is produced in such collisions, it will evidently consist of quarks and gluons – the crucial feature to check is whether they are deconfined. This can be done indirectly, using the equation of state of the quark-gluon plasma together with a specific evolution pattern as input and then looking for the consequences this leads to for the hadrons emitted at the end of the collision process. On the other hand, we can try to check directly if the quarks and gluons in the primordial state are confined or not. This, however, requires hard probes, capable of resolving the short-distance structure of the quark-gluon plasma; it cannot be achieved by studying the usual large hadrons formed at freeze-out. In this section we want to use J/ψ production to show how confined and deconfined matter can be distinguished by hard probes [15].

For matter in a confined state, the constituents will be hadrons; for simplicity, consider an ideal gas of pions. Their momentum distribution will be thermal, i.e., for temperatures not too low it will be given by $\exp(-E_\pi/T) \simeq \exp(-p_\pi/T)$. Hence the average momentum of a pion in this medium is $\langle p_\pi \rangle = 3T$. The distribution of quarks and gluons within a pion is known from structure function studies; the gluon density is $g(x) \simeq 0.5(1-x)^3$ for large $x = p_g/p_\pi$.² As a consequence, the average momentum of a gluon in confined matter is given by

$$\langle p_g \rangle_{\text{conf}} = \frac{1}{5} \langle p_\pi \rangle = \frac{3}{5}T. \quad (15)$$

Hence in a medium of temperature $T \simeq 0.2$ GeV, the average gluon momentum is around 0.12 GeV.

In contrast, the distribution of gluons in a deconfined medium is directly thermal, i.e., $\exp(-p_g)$, so that

$$\langle p_g \rangle_{\text{deconf}} = 3T. \quad (16)$$

Hence the average momentum of a gluon in a deconfined medium is five times higher than in a confined medium³; for $T = 0.2$, it becomes 0.6 GeV. We thus have to find a way to look for such a hardening of the gluon distribution in deconfined matter.

The lowest $c\bar{c}$ vector state, the J/ψ , provides an ideal probe for this. It is very small, with a radius $r_\psi \simeq 0.2$ fm $\ll \Lambda_{QCD}^{-1}$, so that J/ψ interactions with the conventional light quark hadrons probe the short distance features, the parton infrastructure, of the latter. It is very strongly bound, with a binding energy $\epsilon_\psi \simeq 0.65$ GeV $\gg \Lambda_{QCD}$; hence it can be broken up only by hard partons. Since it shares no quarks or antiquarks with pions or nucleons, the dominant perturbative interaction for such a break-up is the exchange of a hard gluon.

We thus see qualitatively how a deconfinement test can be carried out. If we put a J/ψ into matter of temperature $T = 0.2$ GeV, then

- if the matter is confined, $\langle p_g \rangle_{\text{conf}} \simeq 0.12$ GeV, which is much too soft to resolve the J/ψ as $c\bar{c}$ as a bound $c\bar{c}$ state and much less than the binding energy ϵ_ψ , so that the J/ψ survives;

² For very small x , recent results from HERA indicate a steeper increase; however, this does not effect our considerations here.

³ We could equally well assume matter at a fixed energy density, instead of temperature. This would lead to gluons which in case of deconfinement are approximately three times harder than for confinement.

- if the matter is deconfined, $\langle p_g \rangle_{\text{decon}} \simeq 0.6$ GeV, which with some momentum spread is hard enough to resolve the J/ψ and to break the binding, so that the J/ψ will disappear.

Our arguments thus provide a dynamical basis for J/ψ suppression by colour screening [16], and they indicate in fact that J/ψ suppression in dense matter will occur if and only if there is deconfinement.

To put these arguments on a firm theoretical basis, we need the cross section for the inelastic interaction of a J/ψ with a light hadron h , $\sigma_{\psi h}$, from which we can then calculate if a J/ψ can be broken up on its passage through hadronic matter. Because of the small radius and large binding energy of the J/ψ , $\sigma_{\psi h}$ can be calculated in the short-distance analysis of QCD [17,18]. The crucial feature for this calculation is the fact that heavy and tightly bound quarkonium states can be broken up by scattering on usual light hadrons only through the exchange of hard gluons. The momentum distribution of gluons within such a light hadron incident on a J/ψ , $g(x)$, with $x = p_g/p_h$, is a universal non-perturbative input, determined e.g. from parton counting rules or from deep inelastic processes. It has for large x in general the form $g(x) \simeq (1-x)^k$, with $k \simeq 3$ for mesons, $k \simeq 4$ for nucleons. The resulting $J/\psi - h$ cross section then becomes [15]

$$\sigma_{\psi h}(s) \simeq 3 \text{ mb} (1 - \lambda_o/E_h)^{k+5/2}, \quad (17)$$

with $E_h = (m_h^2 + p_h^2)^{1/2}$ and $\lambda_o \simeq (\epsilon_\psi + m_h)$. For low collision energies $\sqrt{s}$, the cross section is thus determined by the behaviour of the gluon distribution at large x , as already noted above, and this leads to a very slow growth from threshold towards the asymptotic value of 3 mb. The functional form of this behaviour is shown in Fig. 5.

Figure 5. Energy dependence of the J/ψ -nucleon cross-section, from [15].

Only in the limit of large quark mass can the quarkonium-hadron cross section can be calculated rigorously in short-distance QCD. We thus have to ask if the charm quark mass is sufficiently large to apply the results of heavy quark theory. For large

M_Q , the quarkonium binding becomes completely Coulombic, and for Coulomb bound states, radius and binding energy are related by

$$\epsilon_{(Q\bar{Q})} = \frac{1}{M_Q r_{(Q\bar{Q})}^2}. \quad (18)$$

For the J/ψ parameters $\epsilon_\psi \simeq 0.65$ GeV, $r_\psi \simeq 0.2$ GeV and $M_c \simeq 1.5$ GeV, relation (18) is in fact quite well satisfied. For an empirical test, one can use the same approach to calculate the cross section for the photoproduction of open charm, $\sigma_{\gamma h \rightarrow c\bar{c}}(s)$. The result is

$$\sigma_{\gamma h \rightarrow c\bar{c}}(s) \simeq 1.2 \mu\text{b} (1 - \nu_o/\nu)^k, \quad (19)$$

with $2\nu = s - m_h^2$ and $2\nu_o = M_\psi^2 + 2M_\psi m_h$. For this process, there are data [19], and in Fig. 6 it is seen that they agree well with the heavy quark theory result (19).

Figure 6. Energy dependence of open charm photoproduction, from [15].

A direct test of eq. 17 will in fact also soon become possible. The ideal way to study the fate of J/ψ 's in confined matter is in principle provided by hadron-nucleus collisions,

$$h A \rightarrow J/\psi + \dots, \quad (20)$$

where the J/ψ is then observed by measuring its decay dimuons, $J/\psi \rightarrow \mu^+ \mu^-$. In practice, however, the muons are identified by passing through an absorber, and for this they have to be sufficiently energetic. As a consequence, experiments in which a hadron beam is incident on a nuclear target lead to J/ψ 's which are very fast ($x_F \geq 0$) in the rest frame of the nucleus and hence leave the nuclear medium long before ever becoming a fully formed physical resonances. Such experiments thus mainly study the behaviour of a coloured $c\bar{c}$ pair in the medium, and the behaviour of this system is quite different from that of a colour-singlet J/ψ . To test the interaction of the latter in the nuclear medium, we need J/ψ 's in the momentum range $-1 \leq x_F \lesssim -0.3$, and as noted, these lead to dimuons too slow to be observable in the usual pA experiments.

However, with the advent of the Pb-beam at the CERN-SPS, it will become possible to study J/ψ -production from a nuclear beam incident on a hydrogen target. Fast J/ψ 's in the lab system will then be slow in the nuclear rest frame and thus pass through the confined nuclear medium as fully formed physical resonances. The short-distance QCD analysis presented here predicts essentially no absorption for this passage, in contrast to a suppression of 25% or more if the asymptotic break-up cross-section is used instead of eq. (17).

To estimate the survival chances of J/ψ 's in confined hadronic matter of higher than standard nuclear density, we consider the survival probability

$$S = e^{-nL\sigma_{\psi h}}. \quad (21)$$

For hadronic matter of nuclear dimension ($L \simeq 5$ fm) but of 10 times higher density ($n \simeq 10 n_0$), we still find essentially no suppression up to temperatures of 500 MeV. Eq. (21) does not take into account the expansion of the system, which would reduce even further any possible chances for a break-up. We thus conclude that there will not be any noticeable J/ψ absorption in confined hadronic matter of the volume and density obtainable in nuclear collisions, and that only deconfined matter can lead to J/ψ suppression. J/ψ production thus provides us indeed with a clear-cut probe to test if a given sample of strongly interacting matter is confined or deconfined.

In closing, we comment briefly on the interpretation of the J/ψ suppression observed in $O - U$ and $S - U$ collisions at CERN [20]. At this time, there remain at least two open questions which prevent us from concluding that this suppression is due to deconfinement. Our present considerations exclude J/ψ absorption in confined matter. The temporal sequence of J/ψ formation and equilibration in nuclear collision is not yet clear, however, and so energetic pre-equilibrium hadrons could have broken up the suppressed J/ψ 's, either as fully formed resonances or in a nascent state. Such break-up processes can be studied theoretically, and it should be possible to determine their effect. Another uncertainty arises from J/ψ production through χ decay, which leads to a sizeable fraction of the overall production rate (30 - 40 %). Since the χ 's are much less tightly bound, their break-up into $D\bar{D}$ pairs is easier, and hence much of the observed suppression could in fact be χ suppression. In $p - Li$ and $\pi^\pm - Li$ interactions at 300 GeV beam momentum [21], the production rates of the different hidden charm final states (J/ψ , χ_1 , χ_2 , ψ') were measured, and if this would be done also in nuclear collisions, the effect of χ suppression could be determined. Thus both theorists and experimentalists have some work left to do before the relation between the observed J/ψ suppression and colour deconfinement is fully clarified.

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