

CERN-TH.7274/94
BI-TP 94/24**QUARKONIUM INTERACTIONS IN HADRONIC MATTER**

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The cross section for the J/ψ and Υ interaction with light hadrons is calculated in short-distance QCD, based on the large heavy quark mass and the resulting large energy gap to open charm or beauty. The low energy form of the cross section is determined by the gluon structure functions at large x ; hence it remains very small until quite high energies. This behaviour is experimentally confirmed by charm photoproduction data. It is shown to exclude J/ψ absorption in confined hadronic matter of the size or density attainable in nuclear collisions; in contrast, the harder gluon spectrum in deconfined matter allows break-up interactions.

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The interaction of heavy quark-antiquark resonances (quarkonia) with the ordinary light hadrons plays an important role both in the dynamics and in the thermodynamics of strong interaction physics. For QCD dynamics, it is important since the small quarkonium size probes the short-distance aspects of the big light hadrons and thus makes a parton-based calculation of the overall cross section possible [1–4]. In QCD thermodynamics, quarkonia can be used as a probe for deconfinement [5], provided their interaction in dense confined matter can be distinguished from that in a quark-gluon plasma.

In this letter we will first recall the QCD analysis of quarkonium interactions with light hadrons. It concludes that the small size of quarkonia, combined with the rather large mass gap to open charm or beauty, strongly inhibits their break-up by low energy collisions with light hadrons [2]; the total quarkonium-hadron cross section attains a constant asymptotic value only at very high energies, compared to corresponding cross sections for the interaction between light hadrons. This slowly rising form of the cross section is derived from an operator product expansion with ensuing sum rules and becomes quite transparent in parton language. It is experimentally well supported by applying the corresponding analysis to data for charm photoproduction [6–9].

We then use this result to show that a slow J/ψ or Υ cannot be absorbed (i.e., broken up into a $D\bar{D}$) in hadronic matter of the volume and density which can be obtained in nuclear collisions. The essential point is here that absorption requires interactions with sufficiently hard gluons, and the gluon distribution within hadrons does not provide these for matter at meaningful temperatures. Hence absorption in confined hadronic matter is excluded as a possible cause of the J/ψ suppression observed in nucleus-nucleus collisions. On the other hand, we find that quarkonia can be broken in deconfined matter, which contains much harder gluons. This provides a microscopic basis of quarkonium suppression by colour screening [5,10].

The QCD analysis of quarkonium interactions applies to heavy and strongly bound quark-antiquark states [2]; therefore we here restrict ourselves to the lowest $c\bar{c}$ and $b\bar{b}$ vector states J/ψ and Υ , which we denote generically by Φ , following the notation of [2]. For such states, both the masses m_Q of the constituent quarks and the binding energies $\epsilon_0(\Phi) \simeq (2M_{(Qq)} - M_\Phi)$ are much larger than the typical scale Λ_{QCD} for non-perturbative interactions; here (Qq) denotes the lowest open charm or beauty state. In $\Phi - h$ interactions, as well as in Φ -photoproduction, $\gamma h \rightarrow \Phi h$, we thus only probe a small spatial region of the light hadron h ; these processes are much like deep-inelastic lepton-hadron scattering, with large m_Q and ϵ_0 in place of the large virtual photon mass $\sqrt{-q^2}$. As a result, the calculation of Φ -photoproduction and of absorptive $\Phi - h$ interactions can be carried out in the short-distance formalism of QCD. Just like deep-inelastic lepton production, these reactions probe the parton structure of the light hadron, and so the corresponding cross sections can be calculated in terms of parton interactions and structure functions.

In the following, we shall first sketch the theoretical basis which allows quarkonium interactions with light hadrons to be treated by the same techniques as used in deep-inelastic lepton-hadron scattering or in the photoproduction of charm. We show the derivation for the sum rules which relate the absorptive $\Phi - h$ cross section to hadronic

gluon structure functions [1,2]. This relation given, we calculate explicitly the energy dependence of the cross section. Readers only interested in this behaviour can therefore go immediately to eq. (24).

Consider the amplitude for forward scattering of a virtual photon on a nucleon,

$$F(s, q^2) \sim i \int d^4x e^{iqx} \langle N | T \{ J_\mu(x) J_\nu(0) \} | N \rangle. \quad (1)$$

In the now standard application of QCD to deep-inelastic scattering one exploits the fact that at large spacelike photon momenta q the amplitude is dominated by small distances of order $1/\sqrt{-q^2}$ (Fig. 1a). The Wilson operator product expansion then allows the evaluation of the amplitude at the unphysical point $pq \rightarrow 0$, where p is the four-momentum of the nucleon. Since the imaginary part of the amplitude (1) is proportional to the experimentally observed structure functions of deep-inelastic scattering, the use of dispersion relations relates the value of the amplitude at $pq \rightarrow 0$ point to the integrals over the structure functions, leading to a set of dispersion sum rules [11]. The parton model can be considered then as a particularly useful approach satisfying these sum rules.

In the case $J_\mu = \bar{Q}\gamma_\mu Q$, i.e., when vector electromagnetic current in eq. (1) is that of a heavy quark-antiquark pair, large momenta q are not needed to justify the use of perturbative methods. Even if $q \sim 0$, the small space-time scale of x is set by the mass of the charmed quark, and the characteristic distances which are important in the correlator (1) are of the order of $1/2m_Q$ (Fig. 1b). In [3,4], this observation was used to derive sum rules for charm photoproduction in a manner quite similar to that used for deep-inelastic scattering.

In the interaction of quarkonium with light hadrons, again the small space scale is set by the mass of the heavy quark, and the characteristic distances involved are of the order of quarkonium size, i.e., smaller than the non-perturbative hadronic scale Λ_{QCD}^{-1} . Moreover, since heavy quarkonium and light hadrons do not have quarks in common, the only allowed exchanges are purely gluonic. However, the smallness of spatial size is not enough to justify the use of perturbative expansion [2]. Unlike in the case of Φ -photoproduction, heavy quark lines now appear in the initial and final states (see Fig. 1c), so that the $(Q\bar{Q})$ state can emit and absorb gluons at points along its world line widely separated in time. These gluons must be hard enough to interact with a compact colour singlet state (colour screening leads to a decoupling of soft gluons with the wavelengths larger than the size of the Φ); however, the interactions among the gluons can be soft and nonperturbative. We thus have to assure that the process is compact also in time. Since the absorption or emission of a gluon turns a colour singlet quarkonium state into a colour octet, the scale which regularizes the time correlation of such processes is by the quantum-mechanical uncertainty principle just the mass difference between the colour-octet and colour-singlet states of quarkonium: $\tau_c \sim 1/(\epsilon_8 - \epsilon_1)$. The perturbative Coulomb-like piece of the heavy quark-antiquark interaction

$$V_k(r) = -g^2 \frac{c_k}{4\pi r} \quad (2)$$

is attractive in the colour singlet ($k = 1$) and repulsive in the colour-octet ($k = 8$) state;

in SU(N) gauge theory

$$V_1 = -\frac{g^2}{8\pi r} \frac{N^2 - 1}{N}, \quad (3a)$$

$$V_8 = \frac{g^2}{8\pi r} \frac{1}{N}. \quad (3b)$$

To leading order in $1/N$, the mass gap between the singlet and octet states is therefore just the binding energy of the heavy quarkonium ϵ_0 , and the characteristic correlation time for gluon absorption and emission is

$$\tau_c \sim 1/\epsilon_0. \quad (4)$$

Although the charm quark is not heavy enough to ensure a pure Coulomb regime even for the lowest $c\bar{c}$ bound states (η_c and J/ψ), the mass gap determined from the observed value of open charm threshold clearly shows that $\tau_c < \Lambda_{QCD}$. For the Υ , the interaction is in fact essentially Coulomb-like and the mass gap to open beauty is even larger than for charm. One therefore expects to be able to treat quarkonium interactions with light hadrons by the same QCD methods that are used in deep-inelastic scattering and charm photoproduction.

We thus use the operator product expansion to compute the amplitude of heavy quarkonium interaction with light hadrons,

$$F_{\Phi h} = i \int d^4x e^{iqx} \langle h | T \{ J(x) J(0) \} | h \rangle = \sum_n c_n(Q, m_Q) \langle O_n \rangle, \quad (5)$$

where the set $\{O_n\}$ includes all local gauge invariant operators expressible in terms of gluon fields; the matrix elements $\langle O_n \rangle$ are taken between the initial and final light-hadron states. The coefficients c_n are computable perturbatively [1] and process-independent. As noted above, in deep-inelastic scattering the expansion (5) is useful only in the vicinity of the point $pq \rightarrow 0$. The same is true for the case of quarkonium interaction with light hadrons. As shown in [2], the expansion (5) can therefore be rewritten as an expansion in the variable

$$\lambda = \frac{pq}{M_\Phi} = \frac{1}{2M_\Phi} (s - M_\Phi^2 - M_h^2), \quad (6)$$

where M_h is the mass of the light hadron. For the lowest 1S quarkonium state one then obtains

$$F_{\Phi h} = r_0^3 \epsilon_0^2 \sum_{n=2}^{\infty} d_n \langle O_n \rangle \left(\frac{\lambda}{\epsilon_0} \right)^n, \quad (7)$$

where r_0 and ϵ_0 are Bohr radius and binding energy of the quarkonium, and the sum runs over even values of n to ensure the crossing symmetry of the amplitude. The most important coefficients d_n were computed in [1] to leading order in g^2 and $1/N$.

Since the total $\Phi - h$ cross section $\sigma_{\Phi h}$ is proportional to the imaginary part of the amplitude $F_{\Phi h}$, the dispersion integral over λ leads to the sum rules

$$\frac{2}{\pi} \int_{M_h}^{\infty} d\lambda \lambda^{-n} \sigma_{\Phi h}(\lambda) = r_0^3 \epsilon_0^2 d_n \langle O_n \rangle \left(\frac{1}{\epsilon_0} \right)^n. \quad (8)$$

Eq. (7) provides only the inelastic intermediate states in the unitarity relation, since direct elastic scattering leads to contributions of order r_0^6 ; hence the total cross section in eq. (8) is due to absorptive interactions only [2]. Recalling now the expressions for radius and binding energy of 1S Coulomb bound states of a heavy quark-antiquark pair,

$$r_0 = \left(\frac{16\pi}{3g^2} \right) \frac{1}{m_Q}, \quad (10a)$$

$$\epsilon_0 = \left(\frac{3g^2}{16\pi} \right)^2 m_Q, \quad (10b)$$

and using the coefficients d_n from [1], it is possible [2] to rewrite these sum rules in the form

$$\int_{M_h}^{\infty} \frac{d\lambda}{M_h} \left(\frac{\lambda}{M_h} \right)^{-n} \sigma_{\Phi h}(\lambda) = 2\pi^{3/2} \left(\frac{16}{3} \right)^2 \frac{\Gamma(n + \frac{5}{2})}{\Gamma(n + 5)} \left(\frac{16\pi}{3g^2} \right) \left(\frac{M_h}{\epsilon_0} \right)^{n-1} \frac{1}{m_Q^2} \langle O_n \rangle. \quad (11)$$

The contents of these sum rules become more transparent in terms of the parton model. In parton language, the expectation values $\langle O_n \rangle$ of the operators composed of gluon fields can be expressed as Mellin transforms [12] of the gluon structure function of the light hadron, evaluated at the scale $Q^2 = \epsilon_0^2$,

$$\langle O_n \rangle = \int_0^1 dx x^{n-2} g(x, Q^2 = \epsilon_0^2). \quad (12)$$

Defining now

$$y = \frac{M_h}{\lambda}, \quad (13)$$

we can reformulate eq. (11) to obtain

$$\int_0^1 dy y^{n-2} \sigma_{\Phi h}(M_h/y) = I(n) \int_0^1 dx x^{n-2} g(x, Q^2 = \epsilon_0^2), \quad (14)$$

with $I(n)$ given by

$$I(n) = 2\pi^{3/2} \left(\frac{16}{3} \right)^2 \frac{\Gamma(n + \frac{5}{2})}{\Gamma(n + 5)} \left(\frac{M_h}{\epsilon_0} \right)^{n-1} \left(\frac{16\pi}{3g^2} \right) \frac{1}{m_Q^2}. \quad (15)$$

Eq. (14) relates the $\Phi - h$ cross section to the gluon structure function. To get a first idea of this relation, we neglect the n -dependence of $I(n)$ compared to that of $\langle O_n \rangle$; then we conclude that

$$\sigma_{\Phi h}(M_h/x) \sim g(x, Q^2 = \epsilon_0^2), \quad (16)$$

since all order Mellin transforms of these quantities are equal up to a constant. From eq. (16) it is clear that the energy dependence of the $\Phi - h$ cross section is entirely determined by the x -dependence of the gluon structure function. The small x behaviour of the structure function governs the high energy form of the cross section, and the hard tail of

the gluon structure function for $x \rightarrow 1$ determines the energy dependence of $\sigma_{\Phi h}$ close to the threshold.

To obtain relation (16), we have neglected the n -dependence of the function $I(n)$. Let us now try to find a more accurate solution of the sum rules (13). We are primarily interested in the energy region not very far from threshold, i.e.,

$$M_h \lesssim \lambda \lesssim 3 \text{ GeV}, \quad (17)$$

since we want to calculate in particular the absorption of Φ 's in confined hadronic matter. In such an environment, the constituents will be hadrons with momenta of at most a GeV or two. A usual hadron (π , ρ , nucleon) of 3 GeV momentum, incident on a J/ψ at rest, leads to $\sqrt{s} \simeq 5.5$ GeV, and this corresponds to $\lambda \simeq 3$ GeV.

From what we learned above, the energy region corresponding to the range (17) will be determined by the gluon structure function at values of x not far from unity. There the x -dependence of $g(x)$ can be well described by a power law

$$g(x) = g_2 (k+1) (1-x)^k, \quad (18)$$

where the function (18) is normalized so that the second moment (12) gives the fraction g_2 of the light hadron momentum carried by gluons, $\langle O_2 \rangle = g_2 \simeq 0.5$. This suggests a solution of the type

$$\sigma_{\Phi h}(y) = a(1-y)^\alpha, \quad (19)$$

where a and α are constants to be determined. Substituting (18) and (19) into the sum rule (13) and performing the integrations, we find

$$a \frac{\Gamma(\alpha+1)}{\Gamma(n+\alpha)} = \left(\frac{2\pi^{3/2} g_2}{m_Q^2} \right) \left(\frac{16}{3} \right)^2 \left(\frac{16\pi}{3g^2} \right) \left(\frac{M_h}{\epsilon_0} \right)^{n-1} \frac{\Gamma(n+\frac{5}{2}) \Gamma(k+2)}{\Gamma(n+5) \Gamma(k+n)}. \quad (20)$$

We are interested in the region of low to moderate energies; this corresponds to relatively large x , to which higher moments are particularly sensitive. Now $\epsilon_0 \simeq M_h$, if we consider nucleons or ρ 's as typical hadrons (the case of pions requires a separate treatment); hence for the range of n for which eq. (5) is valid [4], $n \lesssim 8$, the essential n -dependence is contained in the Γ -functions. For $n \gtrsim 4$, eq. (20) can be solved in closed form by using an appropriate approximation for the Γ -functions. We thus obtain

$$a \frac{\Gamma(\alpha+1)}{\Gamma(k+2)} \simeq \text{const. } n^{\alpha-k-5/2}. \quad (21)$$

Hence to satisfy the sum rules (14), we need

$$\alpha = k + \frac{5}{2} \quad a = \text{const. } \frac{\Gamma(k+2)}{\Gamma(k+\frac{7}{2})}. \quad (22)$$

Therefore the solution of the sum rules (13) for moderate energies λ takes the form*

$$\sigma_{\Phi h}(\lambda) = 2\pi^{3/2} g_2 \left(\frac{16}{3} \right)^2 \left(\frac{16\pi}{3g^2} \right) \frac{1}{m_Q^2} \frac{\Gamma(k+2)}{\Gamma(k+\frac{7}{2})} \left(1 - \frac{M_h}{\lambda} \right)^{k+5/2}. \quad (23)$$

* The same functional form for the energy dependence was also obtained in [2] through parton model arguments.

To be specific, we now consider the J/ψ -nucleon interaction. Setting $k = 4$ in accord with quark counting rules, using $g_2 \simeq 0.5$ and expressing the strong coupling g^2 in terms of the binding energy ϵ_0 (eq. 10b), we then get from (23) the energy dependence of the $J/\psi N$ total cross section

$$\sigma_{J/\psi N}(\lambda) \simeq 2.5 \text{ mb} \times \left(1 - \frac{M_N}{\lambda}\right)^{6.5}, \quad (24)$$

with λ given by eq. (6). This cross section rises very slowly from threshold, as shown in Fig. 2; for $\sqrt{s} = 5 \text{ GeV}$ ($P_N \simeq 2 \text{ GeV}$), it is around 0.1 mb, i.e., more than an order of magnitude below its asymptotic value. This clearly shows that hadronic matter at meaningful densities and for the volumes relevant for nuclear collisions is not efficient in breaking up quarkonium states. To illustrate this, we note that in a volume of the size of a uranium nucleus, the average path length for a J/ψ is $L \simeq (3/4)(1.2 \text{ A}^{1/3}) \simeq 5.6 \text{ fm}$. For matter of standard nuclear density $n_0 = 0.17 \text{ fm}^{-3}$, a cross section $\sigma_{J/\psi N} = 0.1 \text{ mb}$ leads to a survival probability $S = \exp(-\sigma_{J/\psi N} n_0 L) \simeq 0.99$, and even at five times standard nuclear density, we still have a survival probability of more than 95%. – We should note that the high energy cross section of 2.5 mb in eq. (24) is *calculated* in the short-distance formalism of QCD and determined numerically by the values of m_c and ϵ_0 .* It agrees very well with the 2 - 3 mb obtained from photoproduction data via vector meson dominance [6].

For sufficiently heavy quarkonia, the calculation of the $\Phi - h$ cross section as presented is firmly based on short-distance QCD. To check its applicability for charmonium, experimental tests are certainly important. This is indeed possible, since a completely analogous derivation [4] provides the cross section of charm photoproduction, $\gamma N \rightarrow c\bar{c} X$. Instead of eq. (24), we then get

$$\sigma_{\gamma N \rightarrow c\bar{c}}(\nu) \simeq 1.2 \text{ } \mu\text{b} \times \left(1 - \frac{\nu_0}{\nu}\right)^4, \quad (25)$$

in terms of the scaling variable $\nu = (s - M_N^2)/2$, with $\nu_0 = [(M_{J/\psi} + M_N)^2 - M_N^2]/2$. The exponent in eq. (25) is again determined by the gluon structure function; it differs from that in eq. (24) because of the absence of the bound state Coulomb wave function in computing the coefficient function for $c\bar{c}$ photoproduction. In Fig. 3 we compare the result (25) to a summary of data [6–9]. Since eq. (25) contains the contributions of both open and hidden charm, we have reduced it in Fig. 3 by 25% to account for the absence of elastic shadow scattering in the data. It is seen to agree quite well with the experimental behaviour, which in particular shows clearly the small cross section in the threshold region.

A similar test is in principle provided by J/ψ -photoproduction, $\gamma N \rightarrow J/\psi N$. The forward scattering amplitude of this process will give the total charm photoproduction

* An alternative and in principle equivalent approach is to calculate the $\Phi - N$ cross section by perturbative QCD [13]. The sum rule formalism used here has the advantage of providing the total cross section directly in terms of the largely model-independent scales m_Q , ϵ_Q and the fraction g_2 .

cross section $\sigma_{\gamma N \rightarrow c\bar{c}}$, provided vector meson dominance holds and the real part of the amplitude can be neglected. Near threshold this is presumably not the case, and so (Fig. 4) the result of our calculation falls there below the data [6], even though also here the cross section attains its high energy value very slowly.

The calculation of the total Φh cross section as given above includes all possible quarkonium break-up mechanisms and becomes exact in the heavy quark limit. It is interesting, however, to consider also non-perturbative mechanisms made possible by the finiteness of the charmed quark mass. Here the “subthreshold” process [14]

$$J/\Psi + N \rightarrow \Lambda_c + \bar{D} \quad (26)$$

deserves particular attention, since it can occur even at very small energies. Apart from the perturbative contribution calculated above, the amplitude for reaction (26) can acquire a non-perturbative piece, corresponding to the tunneling of the c -quark from the charmonium potential well over a large distance $R_t \sim \Lambda_{QCD}^{-1}$ into the potential well of the nucleon. One can roughly estimate the probability of this process. The time τ_t required for the tunneling over the distance R_t is $\tau_t \sim \Lambda_{QCD}^{-1}$. On the other hand, the energy fluctuation needed to overcome the potential barrier of the height ϵ_0 has according to the uncertainty relation a much shorter lifetime, of order $\tau_f \sim \epsilon_0^{-1}$. The probability for tunneling is thus expected to be

$$W_t \simeq \exp(-\tau_t/\tau_f) \simeq \exp(-\epsilon_0/\Lambda_{QCD}). \quad (27)$$

This probability vanishes as expected in the heavy quark limit, since $\epsilon_0 \sim m_Q$ (see eq. (10b)). For the J/ψ , the probability (27) is already quite small (about 0.02), and hence the corresponding non-perturbative contribution is unlikely to affect the total quarkonium absorption cross section in any significant way.

Since we have concluded that quarkonia cannot be broken up in hadronic matter, we should comment briefly on the implication of our results for hA collisions. In presently available data, the produced quarkonia (J/ψ , ψ' , Υ) have very high momenta in the rest frame of the target nucleus. Hence they are not yet physical resonances when they leave the nuclear medium; as an experimental consequence of this, J/ψ and ψ' , although very different in size, suffer the same nuclear suppression [15]. The observed distributions for these nascent quarkonia are in fact well described by a combination of quantum-mechanical coherence effects (“shadowing”) [16] and colour octet interactions [17] of the $Q\bar{Q}$ system in the nucleus. Thus there is so far no experimental information about the interaction of fully formed physical quarkonia in the confined medium of a heavy nucleus. To obtain such information, quarkonium production has to be measured at large negative x_F values [17]. For a hadron beam incident on a nuclear target, this is difficult, since the slow decay dileptons are hard to measure. The relevant experiment is possible, however, for a nuclear beam incident on a hydrogen target, so that the study of quarkonia in nuclear matter becomes feasible with the advent of the Pb-beam at the CERN-SPS. Our analysis predicts no (i.e., less than 5%) suppression for the production of the basic (1S) J/ψ in the range $-1 \leq x_F \lesssim -0.4$, in contrast to some 25% suppression obtained by using the asymptotic Φh cross section even in the threshold

region [18–20]. Since the mass of the ψ' is just at the open charm threshold, it is much easier to dissociate and hence will suffer much stronger absorption. The same holds to a lesser extent also for the χ states.

In closing, we want to compare quarkonium interactions in confined and in deconfined matter. In confined matter, the crucial gluon densities, $g(x) \sim (1-x)^k$, lead to $\langle x \rangle = 1/(k+2)$, which implies

$$\langle p_g \rangle = \frac{1}{k+2} \langle P_h \rangle \quad (26)$$

for the average momentum of the gluons in terms of that of the hadron incident on the quarkonium. For hadronic matter of temperature $T \simeq 200$ MeV, we thus get

$$\langle p_g \rangle \simeq \frac{3}{5} T \simeq 0.12 \text{ GeV} \quad (27a)$$

for the case of (massless) pions, with $P_\pi \simeq 3T$ and $k = 3$, and

$$\langle p_g \rangle \simeq \frac{1}{3} \sqrt{2mT/\pi} \simeq 0.12 \text{ GeV} \quad (27b)$$

for nucleons, with $P_N \simeq 2\sqrt{2mT/\pi}$ and $k = 4$. It is clear that for gluons of such momenta it is difficult to break up quarkonium, i.e., to overcome the threshold to open charm. In previous absorption studies of charmonium suppression [21], it was suggested that this threshold is overcome in the collision of resonances (ρ , ...) with the J/ψ , by making use of the difference between resonance and pion mass, $M_\rho - 2M_\pi$. Such arguments would be correct for the interaction between light and large hadrons; but in the interaction of the tiny quarkonia with light hadrons, the short-distance character due to the heavy quark mass and the large gap to open charm or beauty makes the gluon distribution within the light hadron the important feature, not its mass.

In deconfined matter, on the other hand, we expect gluons in a medium of temperature $T \simeq 0.2$ GeV to have an average momentum

$$\langle p_g \rangle \simeq 3T \simeq 0.6 \text{ GeV}. \quad (28)$$

The impact of gluons with such a momentum spectrum can easily overcome the 0.7 GeV threshold to open charm and hence break up a J/ψ into D 's. The essential feature of deconfinement in this “microscopic” view of J/ψ suppression in a quark-gluon plasma [5] is thus the hardening of the gluon spectrum.

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Figure Captions:

- Fig. 1: Deep inelastic scattering (a), heavy flavour photoproduction (b) and quarkonium-nucleon interaction (c).
- Fig. 2: The energy dependence of the quarkonium-nucleon cross section, normalised to its asymptotic value; also shown is the dependence on the momentum P_N of a nucleon incident on a J/ψ at rest.
- Fig. 3: The cross section of open charm photoproduction [6–9], compared to the prediction of the short distance QCD analysis.
- Fig. 4: The total cross section for charm photoproduction [6], compared to the prediction of the short distance QCD analysis, with the real part of the amplitude neglected.