

COLOUR DECONFINEMENT AND QUARKONIUM DISSOCIATION

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1. Introduction

The crucial tool in the search for the quark-gluon plasma (QGP) is a probe to test if the strongly interacting medium produced in nuclear collisions consists of confined or deconfined quarks and gluons. In this survey, we will show how quarkonia can be used as such a tool. First of all, this requires an understanding of the dynamics of quarkonium production in the absence of a thermal medium, i.e., in hadron-hadron collisions. The theory of quarkonium production can today be tested on a variety of different states (J/ψ , χ_c , ψ' , Υ , Υ' , Υ''); it is confirmed by data over a vast range of collision energies, up to 1.8 TeV. Given quarkonium production dynamics, we then have to address three distinct problems to obtain a viable probe. What are the effects of confined and of deconfined matter on the production of the different quarkonium states? How can pre-equilibrium phenomena affect the production? We begin with a short summary of the answers which we shall obtain to these questions.

1.1 Preview

To be specific, we consider charmonium states; but the arguments are in general applicable to bottomonium states as well. The theory of charmonium-hadron interactions predicts that the tightly bound charmonium ground state J/ψ cannot be dissociated by hadronic matter at temperatures below 0.5 GeV. The resulting transparency of confined media to J/ψ 's can be tested experimentally in nuclear matter, by studying J/ψ -production with a Pb -beam incident on a hydrogen or deuterium target. In contrast, a quark-gluon plasma contains deconfined gluons, and these are hard enough to break up a physical J/ψ .

Hadron-nucleus collisions also provide the information needed to eliminate possible initial state effects. By comparing different hard processes with and without possible final state interactions, the role of the nuclear medium on the initial state can be clarified and the final state effects can be singled out.

Combining hadron-nucleus and nucleus-nucleus studies to determine the effects of confined and deconfined systems on quarkonium production, we can thus test colour deconfinement. In a nutshell: once initial state effects are removed, confined systems have no effect on the J/ψ , but suppress the ψ' ; once deconfinement sets in, both J/ψ and ψ' are suppressed, but the ψ' more strongly.

1.2 Probing Colour Deconfinement

After this short preview of things to come, we return to the general problem of establishing quark-gluon plasma formation. In high energy nuclear collisions, two beams of partons collide; the partons are initially confined to the colliding nucleons. This confinement can be checked, e.g., by studying primary high mass Drell-Yan dilepton production and observing, except for possible nuclear shadowing effects,

the same parton distribution functions as in deep inelastic lepton-nucleon collisions. After the primary collision, we expect abundant multiple interactions, leading to a rapid increase of entropy, quick thermalisation and hence the production of strongly interacting matter. The fundamental question is whether confinement survives this thermalisation. If it does, we have hadronic matter – if not, a quark-gluon plasma. We expect confinement to be lost if the parton density sufficiently surpasses that present in a hadron-hadron interaction, so that partons can no longer be assigned to specific hadrons. How can we check if this has happened?

The QGP is a *dense* system of *deconfined* quarks and gluons. Its density is in fact the reason for deconfinement: in a sufficiently dense medium, the long-range confining forces become screened, so that only short-range ($\ll \Lambda_{\text{QCD}}^{-1}$) interactions between quarks and gluons remain. To study such a medium and determine its nature, we therefore need probes which are *hard* enough to resolve the short sub-hadronic scales and which can distinguish between *confined* and *deconfined* quarks and gluons. In addition, the probe must survive the subsequent evolution of the medium; therefore it certainly cannot be in equilibrium with the later stages of matter. Two hard, strongly interacting signals produced before equilibration and distinct from the medium have been proposed as probes for confinement/deconfinement:

- *heavy* quark-antiquark resonances (charmonium, bottonium) [1,2], and
- *hard* quarks or gluons (jets) [3,4].

In the following, we will comment only briefly on jets, in order to point out the relation between the two probes, and then concentrate on quarkonium production. Quarkonium and jet production are rather well understood in hadron-hadron collisions, where they are accounted for in terms of perturbative QCD and hadronic parton distribution functions [5,6]. In both cases, the initially formed state ($Q\bar{Q}$, q or g) is in general coloured, and it has an intrinsic mass or momentum scale much larger than Λ_{QCD} . For jets, this is also the state to be used as probe, since the behaviour of a fast colour charge passing through confined matter differs from that in a deconfined medium [7]. In confined matter, the colour charge loses energy as it passes from one hadron to the next through the “interhadronic” vacuum, and the energy loss is determined by the string tension σ acting on the colour charge as it leaves the field of a hadron [8]. In a hot deconfined medium, the crucial quantities are the colour screening radius and the mean free path; these determine with how many other charges the passing colour charge interacts and how much energy it loses per unit length [9,10]. The fate of a colour charge in the transition region between these two limits is still quite uncertain.

For fast quarkonia, the situation is similar; they will pass through the medium while still in a coloured state [8], and hence they can be used as probe in the same way as jets. In addition, however, we can consider slow quarkonia, which have be-

come full physical resonances within a hadronic volume around the $Q\bar{Q}$ formation point and thus traverse the medium as colour singlets. Since the intrinsic spatial scales of J/ψ and Υ , determined by the heavy quark masses and the binding energies, nevertheless remain much smaller than the hadronic size $\Lambda_{\text{QCD}}^{-1}$, they interact only with the partons within a big, light hadron and not with the hadron as a whole. They are thus able to probe the partonic state of any medium. In particular, they are essentially unaffected by the soft gluons which make up confined matter, while the hard gluons present in a QGP will dissociate them [2].

For both quarkonia and jets, thermal production in the expected temperature range ($T \lesssim 0.5$ GeV) is excluded by the mass or momentum scales involved; we can therefore be quite sure that such signals were produced prior to QGP formation. They will also not reach an equilibrium with later stages of the medium. Hard jets and fast quarkonia require too much of an energy loss for this, while slow quarkonia, as noted, are either dissociated or not affected by the medium.

For both proposed probes, initial state nuclear effects can occur before QGP formation. Primary quarks and gluons may undergo multiple scattering or experience shadowing in the nucleus *before* they interact to form a $Q\bar{Q}$ pair or a hard transverse parton. These effects have to be understood and taken into account before any QGP analysis [11,12]. It is therefore necessary to study them in processes which are not effected by the subsequent medium, such as the production of hard direct photons [13] or of high mass Drell-Yan dileptons [14]. In these cases, we have only annihilation or bremsstrahlung of the incident partons; the resulting electromagnetic signal leaves the system unaffected by any subsequent medium and its evolution. If such reactions show nuclear effects, they are presumably due to initial state phenomena.

After these more general remarks, we will now consider in detail the use of quarkonium production as a probe for deconfinement in dense strongly interacting matter. We concentrate on quarkonium for several reasons. J/ψ suppression was predicted [1] to be the consequence of QGP formation, and such a suppression was subsequently indeed observed in high energy nuclear collisions at the CERN-SPS [15]. This triggered an intensive study of possible alternative origins of such a suppression. Hence the analysis necessary to establish an unambiguous probe for deconfinement has been carried much further here than for jets and can provide a good illustration of what needs to be done before drawing any conclusions. In particular, as noted at the beginning, we must understand theoretically and experimentally the dynamics of the process to be used as probe and

- how it is influenced by initial state nuclear effects,
- how it reacts to confined matter,
- how it reacts to deconfined matter, and
- how it reacts to non-equilibrium systems.

In section 2, we will therefore outline the theory of quarkonium production and compare it to present data [6]. Section 3 will deal with quarkonium production in hadron-nucleus collisions. In particular, we will see how the production of *fast* quarkonia provides us with information on the energy loss of a colour charge in confined matter and on gluon shadowing. Next, in section 4, we will discuss the conceptual basis of quarkonia as deconfinement probe and present the results of the heavy quark theory for quarkonium-hadron interactions. We then show how to test experimentally the resulting transparency of confined matter to slow J/ψ 's. Section 5 will bring a comprehensive comparison of the effect of confined and deconfined media on quarkonium production; here we will recover the melting of the J/ψ in a QGP [1,16] on a microscopic level. In Section 6, we remove initial state nuclear modifications and study the effect of pre-equilibrium deconfinement. Finally, in section 7, we give a a brief summary and an assessment of what we have learned from present data.

2. Quarkonium Production in Hadron-Hadron Collisions

In this section, we shall sketch the basic dynamics of quarkonium production in hadron-hadron collisions; for a summary, see e.g. [6]. We shall speak about charmonium states; but everything said also holds for bottonium.

2.1 Colour Evaporation

The first stage of charmonium formation is the production of a $c\bar{c}$ pair; because of the large quark mass, this process can be described by perturbative QCD (Fig. 2.1). A parton from the projectile interacts with one from the target; the parton distributions within the hadrons are determined e.g. by deep inelastic lepton-hadron scattering. Initially, the $c\bar{c}$ pair will in general be in a colour octet state. It subsequently neutralises its colour and binds to a physical resonance, such as J/ψ , χ_c or ψ' . Colour neutralisation occurs by interaction with the surrounding colour field; this and the subsequent resonance binding are presumably of non-perturbative nature (“colour evaporation” [17]). In the evaporation process, the $c\bar{c}$ pair can either combine with light quarks to form open charm mesons (D and \bar{D}) or bind with each other to form a charmonium state.

The basic quantity in this description is the total sub-threshold charm cross section, obtained by integrating the perturbative $c\bar{c}$ production over the mass interval from $2m_c$ to $2m_D$. At high energy, the dominant part of $\tilde{\sigma}_{c\bar{c}}$ comes from gluon fusion (Fig. 2.1a), so that we have

$$\tilde{\sigma}_{c\bar{c}}(s) = \int_{2m_c}^{2m_D} d\hat{s} \int dx_1 dx_2 g_p(x_1) g_t(x_2) \sigma(\hat{s}) \delta(\hat{s} - x_1 x_2 s), \quad (2.1)$$

with $g_p(x)$ and $g_t(x)$ denoting the gluon densities in projectile and target, respectively, and σ the $gg \rightarrow c\bar{c}$ cross section. In pion-nucleon collisions, there are also significant quark-antiquark contributions (Fig. 2.1b), which become dominant at low energies. The essential prediction of the colour evaporation model is that the production cross section of any charmonium state i is given by

$$\sigma_i(s) = f_i \tilde{\sigma}_{c\bar{c}}(s), \quad (2.2)$$

where f_i is a constant which for the time being has to be determined empirically. In other words, the energy dependence of any charmonium production cross section is predicted to be that of the perturbatively calculated sub-threshold charm cross section. As a consequence, the production ratios of different charmonium states

$$\frac{\sigma_i(s)}{\sigma_j(s)} = \frac{f_i}{f_j} = \text{const.} \quad (2.3)$$

are predicted to be energy-independent. – We note that in the generalised colour evaporation model [6], only a part of the total subthreshold cross section $\tilde{\sigma}_{c\bar{c}}$ goes into charmonium formation. In accord with perturbative open charm calculations, the remainder (more than 50 %) leads to $D\bar{D}$ production, with the missing energy obtained by interaction with the colour field.

Fig. 2.1: Production of a $c\bar{c}$ pair by gluon fusion (a) and $q\bar{q}$ annihilation (b)

2.2 Quarkonium Production: Theory and Data

The predictions of the colour evaporation model have recently been compared in a comprehensive survey [6] to the available data, using parton distribution functions [18,19] which take into account the new HERA results. In Figs. 2.2 and 2.3, we see that the energy-dependence is well described for both J/ψ and Υ production; for J/ψ production, the normalisation coefficient is $f_{J/\psi}=0.025$. The Υ results are

obtained for the sum of Υ , Υ' and Υ'' decaying into dimuons, with $Bf_{\Upsilon} = 1.6 \times 10^{-3}$ for the normalisation; here the branching ratios cannot be directly removed. In the fixed target/ISR energy range, the results from the two different sets of parton distributions coincide; for the J/ψ at LHC energies, there is some spread due to scale uncertainties in the parton distributions, which hopefully can be removed by more precise DIS data. For the Υ production, we have already now data up to 1.8 TeV, and in Fig. 2.3 they are seen to agree very well with the prediction obtained using the “low energy” value $Bf_{\Upsilon} = 1.6 \times 10^{-3}$. Previous phenomenological fits [20] had led to much smaller rates; they are included (CR) in Figs. 2.2 and 2.3.

In Figs. 2.4 and 2.5, the predicted energy-independence of production ratios is found to hold as well, again up to Tevatron energy. Here it should be noted that the CDF data for the ratio $\psi'/(J/\psi)$ are taken at large transverse momenta ($5 \leq p_T \leq 15$ GeV), while the lower energy data are integrated over p_T , with the low p_T region dominant. Hence colour evaporation appears to proceed in the same way at both small and large p_T .

The colour evaporation model does not determine the relative production rates of the different states. In order to obtain them, the colour evaporation process has to be specified in more detail. As an example, we consider the ratios of the different $l = 0$ states shown in Figs. 2.4 and 2.5. Assume that the initial colour octet state first neutralises its colour by interaction with the surrounding colour field, producing a colour singlet $c\bar{c}$ state. The relative weights for J/ψ and ψ' production can then be expressed [21] in terms of the corresponding masses and the squared charmonium wave functions at the origin,

$$\frac{\sigma(\psi')}{\sigma(\psi)} = \frac{R_{\psi'}^2(0)}{R_{\psi}^2(0)} \left(\frac{M_{\psi}}{M_{\psi'}} \right)^5. \quad (2.4)$$

Here ψ denotes the directly produced 1S $c\bar{c}$ state, in contrast to the experimentally observed J/ψ , 40% of which originates from radiative χ_c decays [22,23]. The wave functions at the origin can in turn be related to the dilepton decay widths $\Gamma_{ee} \sim (R^2(0)/M^2)$ [24], giving

$$\frac{\sigma(\psi')}{\sigma(\psi)} = \frac{\Gamma_{\psi'}}{\Gamma_{\psi}} \left(\frac{M_{\psi}}{M_{\psi'}} \right)^3. \quad (2.5)$$

Inserting the measured values for masses and decay widths, we find

$$\frac{\sigma(\psi')}{\sigma(\psi)} \simeq 0.24. \quad (2.6)$$

To compare this to the measured value of $(\sigma(\psi')/\sigma(\psi))$, we have to remove the χ_c contributions from the experimental ratio,

$$\frac{\sigma(\psi')}{\sigma(\psi)} = \left[\frac{1}{1 - (\sigma_{\chi}/\sigma_j)} \right] \left[\frac{\sigma(\psi')}{\sigma(J/\psi)} \right]_{\text{exp}}. \quad (2.7)$$

With the experimental values $\sigma(\psi')/\sigma(J/\psi) \simeq 0.14$ (see Figs. 2.2 and 2.3) and $(\sigma_\chi/\sigma_j) \simeq 0.4$ [22,23], this yields $\sigma(\psi')/\sigma(\psi) \simeq 0.23$, in good agreement with the theoretical result (2.6). We thus find that the projection of the colour singlet $c\bar{c}$ state onto J/ψ and ψ' correctly describes their production ratios at all energies and transverse momenta.

The predictions for direct bottomonium production ratios corresponding to Eq. (2.6) are

$$\frac{\sigma(\Upsilon')}{\sigma(\Upsilon)} \simeq 0.36 \quad ; \quad \frac{\sigma(\Upsilon'')}{\sigma(\Upsilon)} \simeq 0.27 . \quad (2.8)$$

Since the contributions from indirect production through radiative χ_b decay are not yet known and there is also feeddown from higher S-states, a quantitative comparison is not possible here. Nevertheless, the predicted values differ by less than 50 % from the data and hence appear reasonable.

3. Quarkonium Evolution and Hadron-Nucleus Collisions

Quarkonium production in a (finite) medium depends crucially on the evolution stage in which the $Q\bar{Q}$ pair is during its passage. Here we determine the different stages attainable in hadron-nucleus collisions and discuss the effects which the nuclear environment has on them.

3.1 Parton Fusion and Colour Neutralisation in Nuclear Matter

The first stage of charmonium production by hadronic interactions is, as we saw, parton fusion resulting in a generally coloured $c\bar{c}$ pair. Subsequently, the $c\bar{c}$ becomes colour neutral by emission or absorption of gluons in the colour field of the interaction, and eventually this small colour singlet state “expands” to form a physical charmonium resonance. If all or part of this evolution takes place inside a nucleus, a number of effects can modify the production:

- the effective distribution of the partons which fuse to form the $c\bar{c}$ can be modified by the nuclear environment (EMC effect);
- the colour octet $c\bar{c}$ state will interact with the nuclear medium strongly and without knowledge of its final charmonium state (nuclear shadowing, $c\bar{c}$ energy loss);
- the physical charmonium states will interact with the nuclear medium according to their different cross sections on hadrons (absorption).

Since the evolution is temporal, with finite time scales, fast and slow charmonia (in the nuclear rest frame) will have quite different fates. In this section, we shall determine the kinematic regimes for the different evolution stages [8,25]; in the subsequent sections, we will then discuss the dominant effects in each stage.

To estimate how long the colour octet $c\bar{c}$ state will live, we first note that the J/ψ wave function keeps the charmed quarks close to their mass shells; we thus have

$p^2 \simeq m_c^2$ (see Fig. 3.1 for notation). The intermediate quark with four-momentum $k = p + q$ is off-shell by an amount

$$\Delta = (p + q)^2 - m_c^2 = 2pq, \quad (3.1)$$

where q is the momentum of the third (colour neutralising) gluon; in the spirit of the parton model, we keep all gluons on-shell. In the low p_T region, in which we are here primarily interested, this third gluon can be arbitrarily soft; colour neutralisation could even involve several gluons. In any case, the colour neutralisation process cannot really be calculated in a purely perturbative framework. We shall here nevertheless keep the structure of Fig. 3.1, taking it to be a phenomenological extension into the non-perturbative soft-gluon regime. A justification of such a procedure might come from recently developed resummation methods [26].

Fig. 3.1: J/ψ production through gluon fusion

The proper life-time of the virtual coloured state is now given by the uncertainty principle as

$$\tau_8 \simeq |\Delta|^{-1/2}. \quad (3.2)$$

In the rest frame of the J/ψ , the life-time of the colour octet,

$$\tau_8 \simeq (2m_c q_0)^{-1/2}, \quad (3.3)$$

is determined by the energy q_0 of the third gluon. It is clear from Eqs. (3.1) and (3.2) that the colour octet state can propagate over long distances only if the third gluon is soft enough. In a confined medium, gluons cannot propagate over distances larger than about 1 fm; hence the low energy cut-off is $q_0 \simeq \Lambda_{\text{QCD}} \simeq 0.2$ GeV. This leads to

$$\tau_8 \simeq 0.25 \text{ fm} \quad (3.4)$$

for the colour neutralisation time. In the rest frame of the nuclear target A , the $c\bar{c}$ travels in this time a distance

$$d_8 = \left(\frac{P_A}{M} \right) \tau_8, \quad (3.5)$$

where P_A is the (lab) momentum of the $c\bar{c}$ and M its mass. From Eq. (3.5) it is clear that in hadron-nucleus collisions sufficiently fast $c\bar{c}$ pairs will still be coloured when they leave the nuclear medium.

The average path length for a $c\bar{c}$ produced in a $h - A$ collision is $(3/4)R_A \simeq 0.86A^{1/3}$ fm; for heavy nuclei, such as Pb ($A = 208$), this becomes about 5 fm. Hence charmonium states of lab momenta $P_A \gtrsim 60$ GeV have passed the nucleus as colour octets. To relate this to the kinematic variables generally used in $h - A$ experiments, we transform the lab momentum P_A to the momentum P measured in the center of mass of a hadron-nucleon collision. The Feynman variable x_F is then defined as $x_F \equiv P/P_{max}$, with $P_{max}(s)$ denoting the maximum cms momentum possible for a charmonium state produced in a hadron-nucleon collision of cms energy \sqrt{s} . From this, we obtain [8,25] that for $\sqrt{s} \gtrsim 20$ GeV and $x_F \gtrsim 0$, the measured charmonium states have traversed most of the nucleus as colour octets. This kinematic range covers essentially all high statistics charmonium production data taken in $h - A$ collisions [27,28]. These experiments therefore study the passage of a colour octet $c\bar{c}$ pair through nuclear matter; as such, they do not provide any information on the interaction of fully formed physical J/ψ or ψ' with nucleons.

An immediate consequence of this is that any nuclear effects observed for charmonium production in $h - A$ collisions in the quoted kinematic range must be the same for J/ψ and for ψ' production. This conclusion is indeed well satisfied experimentally (Fig. 3.2).

We have so far considered the colour regime, i.e., the range in which the $c\bar{c}$ pair traverses the entire nucleus as a colour octet. When the distance d_0 is less than 1 - 1.5 fm, the $c\bar{c}$ is colourless when it passes from the nucleon on which it was produced to the “next” nucleon inside the nucleus. At $\sqrt{s} \simeq 20$ GeV, this occurs for $x_F \lesssim -0.2$. At this point, the $c\bar{c}$ is still a generic small colour singlet state and not yet a specific physical resonance. The other extreme to the colour regime is the range in which the $c\bar{c}$ has become a fully formed physical resonance (J/ψ , ψ') before it leaves the environment of the nucleon at which it was formed. Since the different states have different sizes and hence different formation times, the resonance regime will be state-specific, i.e., at a given collision energy, it will cover a larger x_F -region for the J/ψ than for the ψ' . The distance $d_r(i)$ travelled by the $c\bar{c}$ before becoming a physical resonance state i is given by

$$d_r(i) = \left(\frac{P_A}{M} \right) \tau_r(i) \quad (3.6)$$

where $\tau_r(i)$ denotes the resonance formation time. Potential theory [29] leads to the estimates $t_r(J/\psi) \simeq 0.35$ fm and $t_r(\psi') \simeq 1$ fm. For a collision energy $\sqrt{s} \simeq 20$ GeV, this means that a J/ψ has become a fully formed resonance for $x_F \lesssim -0.5$. The ψ' is still nascent even for $x_F = -1$; although it is a colour singlet on its passage through the nucleus if $x_F \lesssim -0.2$, it has not yet reached its full physical size when it leaves the nucleus even for $x_F = -1$.

In the transition regime between these two extremes, the $c\bar{c}$ experiences colour interactions on part of its passage, while it traverses the medium as a small colour singlet on the remaining part. In Table 1, we have summarised the different kinematic ranges for proton beam incident on a Pb target at a cms nucleon-nucleon collision energy $\sqrt{s} = 20$ GeV, for (1S) J/ψ and (2S) ψ' production. In the following sections, we survey the possible nuclear effects which can arise in the different kinematic regimes.

Regime	J/ψ	ψ'
Colour Regime	$x_F \geq 0$	$x_F \geq 0$
Transition Regime	$-0.5 \leq x_F \leq 0$	$-1 \leq x_F \leq 0$
Resonance Regime	$-1 \leq x_F \leq -0.5$	—

Table 1:
Charmonium Production in $p - A$ Collisions at $\sqrt{s} = 20$ GeV

3.2 Quantum-Mechanical Interference and Nuclear Shadowing

We now address nuclear modifications of charmonium production due to what at first sight appears to be a change in the effective parton distribution function $g_t(x_t)$ in nuclei, compared to that in nucleons. The fraction of the target momentum carried by the corresponding parton in the elementary interactions shown in Fig. 2.1 is given by

$$x_t = \frac{1}{2}(\sqrt{(4M^2/s) + x_F^2} \pm x_F), \quad (3.7)$$

where M is the mass of the $c\bar{c}$ system; the \pm signs hold for $x_F > 0$ and $x_F < 0$, respectively. For $c\bar{c}$ production at $x_F \geq 0$ and $\sqrt{s} \geq 20$, $x_t \leq 0.15$. In this x_t

region, deep inelastic scattering experiments on nuclei find a lower quark density than on nucleons (nuclear shadowing [30]), as illustrated in Fig. 3.3 (the solid line is phenomenological fit [31]). Such nuclear shadowing can, however, not be interpreted as an intrinsic change of the parton distribution in nuclei. As such, it would be applicable in factorisable form to all hard processes; but for Drell-Yan dilepton production by $p - A$ collisions, which involves nuclear quark distribution in the same x_t region, little or no modification is observed [28]. The apparent puzzle is resolved by noting that nuclear shadowing is due to quantum-mechanical interference, similar to the well-known Landau-Pomeranchuk effect [32]. This interference depends on the specific process and hence leads to different nuclear modifications for charmonium hadro- and photoproduction, for Drell-Yan dilepton production and for deep inelastic lepton-hadron scattering [12]. For a first discussion of quantum interference effects in deep inelastic scattering, see [33].

We had seen above in Eq. (3.5) that sufficiently fast $c\bar{c}$ pairs traverse the entire nucleus in a virtual coloured state. In this case, the interactions of the pair with different nucleons cannot be added incoherently and factorisation breaks down even for dynamically uncorrelated nucleons: quantum mechanical interference can now lead to nuclear modifications even though both the parton distribution and the elementary parton interaction amplitude are unchanged. Such interference effects set in when the coherence length $d_c = (1/mx_t)$, over which the $c\bar{c}$ is in an off-shell virtual state, becomes larger than the internucleonic distance $d_0 = n_0^{-1/3} \simeq 1.8$ fm, with $n_0 = 0.17 \text{ fm}^{-3}$ denoting standard nuclear density and m the nucleon mass. This distance has to be compared to the mean free path λ_8 of the virtual colour state in nuclear matter,

$$\lambda_8 = (\sigma_8 n_0)^{-1}, \quad (3.8)$$

where σ_8 is the cross section for the interaction of the colour octet $c\bar{c}$ with nucleons. To estimate the size of σ_8 , we assume that the overall colour neutrality required for the propagation of the coloured $c\bar{c}$ through the nucleus is provided by a comoving light $q\bar{q}$ pair. The interaction cross section of the $Q\bar{Q}q\bar{q}$ system is then defined by the light $q\bar{q}$ pair, so that its size can be as large as that of light $q\bar{q}$ mesons (20 - 30 mb). This is much larger than that of a physical J/ψ ; it is now not the size of the $c\bar{c}$, but its colour which determines the interaction. The mean free path corresponding to such a cross section becomes $\lambda_8 \simeq 2 - 3$ fm, which is about the size of the internucleonic distance $d_0 = n_0^{-1/3} \simeq 1.8$ fm, so that charmonium production in $p - A$ collisions for $x_F \geq 0$ and $\sqrt{s} \geq 20$ GeV satisfies the “shadowing relation” $d_0 \simeq \lambda_8 \ll d_c$ [12]. In this regime, the interactions of the virtual $c\bar{c}$ with successive nucleons interfere destructively with each other; it is the regime which for electromagnetic interactions results in the Landau-Pomeranchuk effect. The production cross section $\sigma(pA \rightarrow c\bar{c})$ on nuclei is here therefore less than the

incoherent result $A\sigma(pp \rightarrow c\bar{c})$.

The basic quantity studied experimentally is the ratio

$$R_{A/p} = \frac{\sigma(pA \rightarrow c\bar{c})}{A\sigma(pp \rightarrow c\bar{c})}; \quad (3.9)$$

it would be unity in the absence of any nuclear effects. It is observed, however, that $R_{A/p}$ decreases with increasing $c\bar{c}$ momentum in the nuclear rest frame (i.e., whenever d_c increases), as well as with increasing A . Rather than in terms of d_c , $R_{A/p}$ is generally studied in terms of the fractional target parton momentum $x_t = (md_c)^{-1}$. Decreasing x_t thus increases the path length for the colour octet inside the nucleus, and a longer path of the virtual state leads to more destructive interference. When d_c reaches the nuclear diameter $2R_A$, we have the maximum possible interference, so that the suppression saturates for $x_t \leq x_t^{\text{sat}} \simeq (2mR_A)^{-1}$, with the approximate limit [12]

$$R_{A/p}^{sh}(x_t^{\text{sat}}) \simeq \frac{\lambda_8}{R_A}. \quad (3.10)$$

A more realistic form including geometric effects is given by

$$R_{A/p}^{sh}(x_t^{\text{sat}}) \simeq 1 - e^{-\lambda_8/R_A}, \quad (3.11)$$

which reduces to Eq. (3.10) for $\lambda/R_A \ll 1$. The available data fall into the region $0 \leq R_A/\lambda_8 = n_0\sigma_8 R_A \lesssim 0.7$, so that the measured nuclear reduction lies between 0.5 and 1. In this region, the relation $e^{-x} \simeq 1 - e^{1/x}$ is quite well satisfied, so that

$$R_{A/p}^{sh}(x_t^{\text{sat}}) \simeq e^{-n_0\sigma_8 R_A} \quad (3.12)$$

also should also describe the data. The form (3.12) was in fact obtained as an empirical fit [34]; we see here, however, that it is the result of quantum-mechanical interference effects of the virtual $c\bar{c}$ and is not related to the interaction of a physical J/ψ with nucleons, as assumed there.

For $d_c \leq d_0$, the scattering on different nucleons becomes incoherent, giving $R_{A/p}^{sh} = 1$, with $x_t \geq x_t^{\text{inc}} \simeq (md_0)^{-1}$. For $A = 200$, the two limits are $x_t^{\text{sat}} = 0.015$ and $x_t^{\text{inc}} = 0.12$. Qualitatively, this is the type of behaviour found in charmonium production for x_F not too large [11]. The specific functional form of the suppression between the two extremes can only be estimated phenomenologically. Experimental results have often been parametrised as $R_{A/p} \simeq A^{\alpha-1}$; our interference considerations imply that the coefficient α depends on x_t . An analysis of charmonium production in $p - A$ collisions [11] suggests a simple linear fit in $\ln A$ and $\ln x_t$,

$$R_{A/p}^{sh} \simeq [1 - c \ln(x_t m d_0) \ln A]. \quad (3.13)$$

The scale $md_0 \simeq 0.12$ is chosen to make $R_{A/p}^{sh} = 1$ in the incoherence regime; the constant c has to be chosen such as to give the correction saturation value. For a fit of $p - A$ data to a similar form, see [8].

3.3 Momentum Loss of a Colour Charge in Nuclear Matter

The suppression $R_{A/p}^{sh}$ cannot be directly compared to data, however, since shadowing is not the only effect suffered by the coloured $c\bar{c}$ state on its path through the nucleus. The $c\bar{c}$ is formed in the collision of the projectile with one of the target nucleons. To leave the nuclear medium, it has to traverse the remaining part of the nucleus, in which it can interact with other nucleons. Getting from one nucleon to the next means passing regions of internuclear physical vacuum, which does not support colour charges. To estimate the effect of this passage, we imagine the $c\bar{c}$ to stretch a string from its formation region to the next nearest nucleon, then from this to the one after that, and so on. If the pair initially had a cms momentum P , then this will be shifted to

$$P' = P - \kappa_8(L_A - d_0), \quad (3.14)$$

where $\kappa_8 \simeq (9/4)\kappa \simeq (9/4)$ GeV/fm is the string tension of the colour octet, $\kappa \simeq 1$ GeV/fm that of the fundamental triplet [8]; L_A is the total path length of the colour state. As consequence of this momentum loss, a $c\bar{c}$ pair observed at a given x_F must have been originally produced at a higher value x_F/δ , with

$$\delta = \left(\frac{P - \kappa_8(L_A - d_0)}{P} \right). \quad (3.15)$$

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

$$F_A(x_F) = W_0 F_0(x_F) + (1 - W_0) \left(\frac{F_0(x_F/\alpha)}{\delta} \right) \Theta(1 - x_F/\delta). \quad (3.16)$$

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

$$W_0 \simeq \exp\{-(L_A - z_0)/d_8\} \quad (3.17)$$

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. (3.16) gives the effect of the n coherent scatterings in the medium, with the resulting shift in x_F . The factor $1/\delta$ assures that the distribution remains normalised. Note that the momentum loss of $c\bar{c}$ 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 . Dividing Eq. (3.16) by the x_F -distribution for h-p collisions, $F_0(x_F)$, gives us the ratio of $h - A$ to $h - p$ distributions,

$$R_{A/p}^{ml}(x_F) = W_0 + (1 - W_0) \left\{ \frac{F_0(x_F/\alpha)}{\alpha F_0(x_F)} \right\} \Theta(1 - x_F/\alpha). \quad (3.18).$$

For a nucleus of $A = 200$ and the other parameters as above, we get the dependence on x_F shown in Fig. 3.4 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 saturates at $\Delta P_{\max} \simeq \kappa_8(R_A - z_0) \simeq 10$ GeV and is thus bounded. As a consequence, the relative shift in x_F decreases with increasing beam energy, and the effect of the medium disappears for $\sqrt{s} \rightarrow \infty$, when $R_{A/p}^{ml} \simeq 1$ for $0 < x_F < 1$.

Fig. 3.4: The effect of colour octet interactions on charmonium production in $p - A$ collisions at three different energies

In addition to the non-perturbative nuclear modification of fast quarkonium production in nuclei, the $c\bar{c}$ pair can undergo hard interactions with the quarks and gluons within each of the nucleons it passes through. We neglect this here, since the density of sufficiently hard scatterings in nuclei is too low to compete with string stretching, as will become evident in Section 3.4. This is changed, however, in a quark-gluon plasma, and the difference in parton hardness for confined and deconfined media is in fact crucial for the use of quarkonia as confinement/deconfinement probe.

The observed suppression of quarkonium production in $h - A$ collisions will be due to both the effects just discussed, nuclear shadowing and momentum loss of

the coloured $c\bar{c}$ pair. We thus have to compare experimental results to the product

$$R_{A/p}(x_F) = R_{A/p}^{sh}(x_F) \times R_{A/p}^{ml}(x_F) \quad (3.19)$$

of the two mechanisms, as given by Eqs. (3.12) and (3.18). The energy loss is effective mainly at larger x_F ; since the integrated cross section, on the other hand, is determined by the region around $x_F \simeq 0$, its suppression is dominantly due to nuclear shadowing. Hence Eq. (3.11) should determine the main A -dependence of charmonium suppression in $p - A$ collisions at $x_f \geq 0$, and as a comparison to the equivalent form (3.12) shows [34], it indeed does.

This form, but with a slightly different parametrisation of shadowing [11], has been compared to the available data [27,28] in [8]. In spite of the simplistic fit to shadowing and some kinematic approximations at small x_F , the result (Fig. 3.5) is seen to reproduce the observed suppression quite well.

Quarkonium production in hadron-nucleus collisions, as experimentally studied up to now, appears thus as theoretically quite well understood, based on the interaction of a fast colour octet $c\bar{c}$ pair with nuclear matter. The dominant mechanisms for this interaction are the destructive quantum-mechanical interference of the scattering amplitudes of the $c\bar{c}$ on different nucleons, and the momentum loss of the $c\bar{c}$ as it traverses the physical vacuum between successive nucleon interactions.

3.4 Charmonium Interactions in Nuclear Matter

If we want to know the effect of nuclear matter on a fully formed physical J/ψ , then – as shown in Table 1 – we have to study the production of J/ψ 's slow in the nuclear rest frame. The study of such slow charmonia has up to now been essentially impossible. It requires the detection of slow decay dileptons, and for this the abundance of slow hadrons constitutes an overwhelming background. In the case of fast dileptons, a hadron absorber can eliminate these, and hence all $p - A$ studies were so far restricted to dilepton pairs of more than 20 GeV in the rest frame of the nuclear target. The advent of the Pb -beam at the CERN-SPS has removed this constraint. With the Pb -beam incident on a hydrogen (or deuterium) target, the nuclear rest frame moves with a lab momentum of 160 GeV. Hence now those charmonia (and their decay dileptons) which are slow in the nuclear rest frame are very fast in the lab system and will thus pass the hadron absorber; such experiments can therefore provide the cross section for the break-up interaction of physical J/ψ 's on nucleons. To avoid confusion, we shall continue to define x_F also in such $Pb - h$ reactions as positive for cms momenta in the direction of the hadron; with this terminology, $Pb - h$ collisions provide us with information about the so far unknown region of negative x_F . Estimates [8,25] show that for $\sqrt{s} \geq 17$ GeV and $x_F \leq -0.2$, a $c\bar{c}$ is colourless before it reaches the next nucleon, so that colour interactions no longer enter.

We now parametrise the survival probability for state i ($i = J/\psi, \chi_c, \psi'$) as

$$S_i = \exp\{-n_0\sigma_i L\}, \quad (3.20)$$

where n_0 is as before normal nuclear density, L the average path length of the charmonium state i in the nucleus, and σ_i its absorption cross section of the state i in nuclear matter. If the charmonium state is fully formed before it leaves the range of the nucleon at which it was produced, σ_i is simply the charmonium-nucleon cross section. If it is not yet fully formed, the effective cross section will be smaller, vanishing in the colour transparency limit of a pointlike colour singlet [35]. For illustration, we can therefore use a simplistic parametrisation of the cross section as function of the distance d the state has travelled before leaving the nucleus [36,37],

$$\sigma_i(d) = \sigma_i \left(\frac{d}{\bar{L}_i} \right)^2, \quad (3.21)$$

where σ_i is the fully developed cross section and $\bar{L} = [(P_A/M_i)\tau_r^{(i)} - 1]$ the effective distance travelled until full resonance formation. Eq. (3.21) holds for $d \leq \bar{L}$; for $d \geq \bar{L}$, $\sigma_i(d) = \sigma_i$. Using this parametrisation, Eq. (3.20) is replaced by

$$S_i = \exp\{-n_0\sigma_i[L - \frac{2}{3}\bar{L}_i]\}; \quad (3.22)$$

this reduces to Eq. (3.20) for momenta high enough to make $\bar{L} = 0$. To determine the actual survival probabilities, we now need the values σ_i for the fully formed resonances colliding with nucleons. That is a topic in its own right and will be taken up in the next chapter. Before turning to this question, we want to connect the survival probability (3.20) to the measured suppression ratio $R_{A/p}$.

3.5 EMC Effect Modifications

The survival probabilities S_i are related to the $p - A$ and $p - p$ induced charmonium production cross sections through

$$\frac{(d\sigma_i^{pA}/dx_F)}{A(d\sigma_i^{pp}/dx_F)} = \frac{g_A(x)}{g_p(x)} S_i(x_F) \equiv R_{A/p}(x) S_i(x_F), \quad (3.23)$$

where $g_A(x)$ and $g_p(x)$ are the parton distribution functions in nuclear and proton target, respectively. The fractional parton momentum x is now, with $x_F \leq 0$, expressed by

$$x = \frac{1}{2}(\sqrt{(x_F^2 + (4M_i^2/s))} - x_F) \quad (3.24)$$

in terms of the variables x_F and s . At the energy possible for Pb -beam experiments ($\sqrt{s} = 17.4$ GeV), and for $x_F \leq 0$, we have $x \geq 0.15$. We are thus above the region

in which quantum-mechanical coherence effects (nuclear shadowing) can play a role; the reaction now is indeed an incoherent sum of interactions with different nucleons. In deep inelastic scattering on nuclear targets there is for $x_2 \geq 0.15$ a modification in comparison to the same process on nucleons (EMC effect [38]; it can now be interpreted as a genuine change of the parton distribution function and applied to other hard processes. The resulting pattern for quark distributions $q(x)$,

$$R_{A/p}^{EMC}(x_F) = \frac{q_A(x_F)}{q_p(x_F)} \quad (3.25)$$

is illustrated in Fig. 3.6 for $M = 3.1$ GeV and $\sqrt{s} = 17.4$ GeV, making use of Eq. (3.24) to convert the Bjorken variable x to x_F .

Although the EMC effect has so far been observed only for quarks, it is to be expected that gluons will exhibit a similar behaviour. This is one reason why the initial state factor $R_{A/p}(x)$ will introduce an further x_F variation in addition to that coming from $S_i(x_F)$. A second reason was already indirectly mentioned above: with increasing $|x_F|$, charmonium production is more and more due to quark-antiquark annihilation rather than to gluon fusion. The two contributions become approximately equal around $|x_F| = 0.5$, and as $|x_F| \rightarrow 1$, the $q\bar{q}$ contribution is dominant. We shall here assume that the gluon distributions behave similarly and use the quark form of $R_{A/p}$ in the whole region $-1.0 \leq x_F \leq 0$. Eqs. (3.22)/(3.23) and the $R_{A/p}^{EMC}$ of Fig. (3.6) then predict the behaviour of the measured cross section ratio in the region of negative x_F .

4. The Theory of Quarkonium-Hadron Interactions

In this section, the cross sections for the inelastic interaction between quarkonia and light hadrons will be calculated in short distance QCD, based on the small radii and the large binding energies of the lowest $Q\bar{Q}$ resonances. A break-up of such states requires hard gluons, but the gluons confined to hadrons slow in the quarkonium rest frame are in general very soft, so that the resulting cross sections remain very small until quite high collision energies.

4.1 The Short-Distance Analysis of Quarkonium-Hadron Interactions

The interaction of quarkonia with 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 [39–43]. In QCD thermodynamics, quarkonia can be used as a probe for deconfinement [1,2], provided their interaction in dense confined matter can be distinguished from that in a quark-gluon plasma.

We begin this chapter with a brief summary of 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 [40]; 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.

The QCD analysis of quarkonium interactions applies to heavy and strongly bound quark-antiquark states [40]; 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 [40]. 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 leptoproduction, 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 [39,40]. 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. (4.24).

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

$$F(s, q^2) \sim i \int d^4x e^{iqx} \langle N | T \{ J_\mu(x) J_\nu(0) \} | N \rangle . \quad (4.1)$$

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. 4.1a). 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 (4.1) 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 [44]. 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. (4.1) 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 (4.1) are of the order of $1/2m_Q$ (Fig. 4.1b). In [41,42], 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 [40]. Unlike in the case of Φ -photoproduction, heavy quark lines now appear in the initial and final states (see Fig. 4.1c), 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

$$V_k(r) = -g^2 \frac{C_k}{4\pi r} \quad (4.2)$$

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

$$V_1 = -\frac{g^2}{8\pi r} \frac{N^2 - 1}{N}, \quad (4.3a)$$

$$V_8 = \frac{g^2}{8\pi r} \frac{1}{N}. \quad (4.3b)$$

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

$$\tau_c \sim 1/\epsilon_0. \quad (4.4)$$

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,

$$F_{\Phi h} = i \int d^4x e^{iqx} \langle h | T \{ J(x) J(0) \} | h \rangle = \sum_n c_n(Q, m_Q) \langle O_n \rangle, \quad (4.5)$$

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 [39] and process-independent. As noted above, in deep-inelastic scattering the expansion (4.5) 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 [40], the expansion (4.5) can therefore be rewritten as an expansion in the variable

$$\lambda = \frac{pq}{M_\Phi} = \frac{(s - M_\Phi^2 - M_h^2)}{2M_\Phi} \simeq \frac{(s - M_\Phi^2)}{2M_\Phi} \quad (4.6)$$

where M_h is the mass of the light hadron; the approximate equality becomes valid the heavy quark limit. For the lowest 1S quarkonium state one then obtains

$$F_{\Phi h} = r_0^3 \epsilon_0^2 \sum_{n=2}^{\infty} d_n \langle O_n \rangle \left(\frac{\lambda}{\epsilon_0} \right)^n, \quad (4.7)$$

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 [39] 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

$$\frac{2}{\pi} \int_{\lambda_0}^{\infty} d\lambda \lambda^{-n} \sigma_{\Phi h}(\lambda) = r_0^3 \epsilon_0^2 d_n \langle O_n \rangle \left(\frac{1}{\epsilon_0} \right)^n. \quad (4.8)$$

Eq. (4.7) 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. (4.8) is due to absorptive interactions only [40], and the integration in Eq. (8) starts at a lower limit $\lambda_0 > M_h$. Recalling now the expressions for radius and binding energy of 1S Coulomb bound states of a heavy quark-antiquark pair,

$$r_0 = \left(\frac{16\pi}{3g^2} \right) \frac{1}{m_Q}, \quad (4.9)$$

$$\epsilon_0 = \left(\frac{3g^2}{16\pi} \right)^2 m_Q, \quad (4.10)$$

and using the coefficients d_n from [39], it is possible [40] to rewrite these sum rules in the form

$$\int_{\lambda_0}^{\infty} \frac{d\lambda}{\lambda_0} \left(\frac{\lambda}{\lambda_0} \right)^{-n} \sigma_{\Phi h}(\lambda) = 2\pi^{3/2} \left(\frac{16}{3} \right)^2 \frac{\Gamma(n + \frac{5}{2})}{\Gamma(n + 5)} \left(\frac{16\pi}{3g^2} \right) \frac{1}{m_Q^2} \langle O_n \rangle, \quad (4.11)$$

with $\lambda_0/\epsilon_0 \simeq 1$ in the heavy quark limit. 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 [45] of the gluon structure function of the light hadron, evaluated at the scale $Q^2 = \epsilon_0^2$,

$$\langle O_n \rangle = \int_0^1 dx x^{n-2} g(x, Q^2 = \epsilon_0^2). \quad (4.12)$$

Defining now

$$y = \frac{\lambda_0}{\lambda}, \quad (4.13)$$

we can reformulate Eq. (11) to obtain

$$\int_0^1 dy y^{n-2} \sigma_{\Phi h}(\lambda_0/y) = I(n) \int_0^1 dx x^{n-2} g(x, Q^2 = \epsilon_0^2), \quad (4.14)$$

with $I(n)$ given by

$$I(n) = 2\pi^{3/2} \left(\frac{16}{3} \right)^2 \frac{\Gamma(n + \frac{5}{2})}{\Gamma(n + 5)} \left(\frac{16\pi}{3g^2} \right) \frac{1}{m_Q^2}. \quad (4.15)$$

Eq. (4.14) 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

$$\sigma_{\Phi h}(\lambda_0/x) \sim g(x, Q^2 = \epsilon_0^2), \quad (4.16)$$

since all order Mellin transforms of these quantities are equal up to a constant. From Eq. (4.16) 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 (4.16), 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 (4.13). We are primarily interested in the energy region not very far from the inelastic threshold, i.e.,

$$(M_h + \epsilon_0) \lesssim \lambda \lesssim 5 \text{ GeV}, \quad (4.17)$$

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 5 GeV momentum, incident on a J/ψ at rest, leads to $\sqrt{s} \simeq 6 \text{ GeV}$, and this corresponds to $\lambda \simeq 5 \text{ GeV}$.

From what we learned above, the energy region corresponding to the range (4.17) 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

$$g(x) = g_2 (k + 1) (1 - x)^k, \quad (4.18)$$

where the function (18) is normalized so that the second moment (4.12) 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

$$\sigma_{\Phi h}(y) = a(1 - y)^\alpha, \quad (4.19)$$

where a and α are constants to be determined. Substituting (4.18) and (4.19) into the sum rule (4.13) and performing the integrations, we find

$$a \frac{\Gamma(\alpha + 1)}{\Gamma(n + \alpha)} = \left(\frac{2\pi^{3/2}g_2}{m_Q^2} \right) \left(\frac{16}{3} \right)^2 \left(\frac{16\pi}{3g^2} \right) \frac{\Gamma(n + \frac{5}{2}) \Gamma(k + 2)}{\Gamma(n + 5) \Gamma(k + n)}. \quad (4.20)$$

We are interested in the region of low to moderate energies; this corresponds to relatively large x , to which higher moments are particularly sensitive. Hence for the range of n for which Eq. (4.5) is valid [42], $n \lesssim 8$, the essential n -dependence is contained in the Γ -functions. For $n \gtrsim 4$, Eq. (4.20) can be solved in closed form by using an appropriate approximation for the Γ -functions. We thus obtain

$$a \frac{\Gamma(\alpha + 1)}{\Gamma(k + 2)} \simeq \text{const. } n^{\alpha - k - 5/2}. \quad (4.21)$$

Hence to satisfy the sum rules (4.14), we need

$$\alpha = k + \frac{5}{2} \quad a = \text{const.} \frac{\Gamma(k+2)}{\Gamma(k+\frac{7}{2})}. \quad (4.22)$$

Therefore the solution of the sum rules (4.13) for moderate energies λ takes the form

$$\sigma_{\Phi h}(\lambda) = 2\pi^{3/2} g_2 \left(\frac{16}{3}\right)^2 \left(\frac{16\pi}{3g^2}\right) \frac{1}{m_Q^2} \frac{\Gamma(k+2)}{\Gamma(k+\frac{7}{2})} \left(1 - \frac{\lambda_0}{\lambda}\right)^{k+5/2}. \quad (4.23)$$

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. 4.10), we then get from Eq. (4.23) the energy dependence of the $J/\psi N$ total cross section

$$\sigma_{J/\psi N}(\lambda) \simeq 2.5 \text{ mb} \times \left(1 - \frac{\lambda_0}{\lambda}\right)^{6.5}, \quad (4.24)$$

with λ given by Eq. (6) and $\lambda_0 \simeq (M_N + \epsilon_0)$. This cross section rises very slowly from threshold, as shown in Fig. 4.2; for $P_N \simeq 5$ GeV, it is around 0.1 mb, i.e., more than an order of magnitude below its asymptotic value.

We should note that the high energy cross section of 2.5 mb in Eq. (4.24) is *calculated* in the short-distance formalism of QCD and determined numerically by the values of m_c and ϵ_0 . From Eqs. (4.10) and (4.23), it is seen to be proportional to $1/(m_Q \sqrt{m_Q \epsilon_0})$. For $\Upsilon - N$ interactions, with $m_b \simeq 4.5$ GeV and $\epsilon_0 \simeq 1.10$ GeV, we thus have the same form (4.24), but with

$$\sigma_{\Upsilon N} \simeq 0.37 \text{ mb} \quad (4.25)$$

as high energy value. This is somewhat smaller than that obtained from geometric arguments [46] and potential theory [47].

For sufficiently heavy quarks, the dissociation of quarkonium states by interaction with light hadrons can thus be fully accounted for by short-distance QCD. Such perturbative calculations become valid when the space and time scales associated with the quarkonium state, r_Q and t_Q , are small in comparison to the nonperturbative scale $\Lambda_{\text{QCD}}^{-1}$

$$r_Q \ll \Lambda_{\text{QCD}}^{-1}, \quad (4.26a)$$

$$t_Q \ll \Lambda_{\text{QCD}}^{-1}; \quad (4.26b)$$

$\Lambda_{\text{QCD}}^{-1}$ is also the characteristic size of the light hadrons. In the heavy quark limit, the quarkonium binding becomes Coulombic, and the spatial size $r_Q \sim (\alpha_s m_Q)^{-1}$

thus is small. The time scale is by the uncertainty relation given as the inverse of the binding energy $E_Q \sim m_Q$ and hence also small. For the charmonium ground state J/ψ , we have

$$r_{J/\psi} \simeq 0.2 \text{ fm} = (1 \text{ GeV})^{-1}; \quad E_{J/\psi} = 2M_D - M_{\psi'} \simeq 0.64 \text{ GeV}. \quad (4.27)$$

With $\Lambda_{\text{QCD}} \simeq 0.2 \text{ GeV}$, the inequalities (4.25) seem already reasonably well satisfied, and also the heavy quark relation $E_{J/\psi} = (1/m_c r_{J/\psi}^2)$ is very well fulfilled. We therefore expect that the dissociation of J/ψ 's in hadronic matter will be governed by the J/ψ -hadron break-up cross section as calculated in short-distance QCD.

Nevertheless, in view of the finite charm quark mass, it makes sense to ask if the formalism developed here correctly describes J/ψ interactions with light hadrons. We will take up this question first theoretically, in the next section, where we shall see if non-perturbative interactions can lead to significant contributions to J/ψ break-up. In Section 4.3, we then see how to verify experimentally the short-distance QCD prediction for J/ψ break-up by light hadrons.

4.2 Non-Perturbative Quarkonium Dissociation

For an isolated J/ψ -hadron system, non-perturbative interactions can be pictured most simply as a quark rearrangement [48]. Consider putting a J/ψ “into” a stationary light hadron; the quarks could then just rearrange their binding pattern to give rise to transitions such as $J/\psi + N \rightarrow \Lambda_c + \bar{D}$ or $J/\psi + \rho \rightarrow D + \bar{D}$ (Fig. 4.3). The probability for such a transition can be written as

$$P_{\text{rearr}} \sim \int d^3r R(r) |\phi_\psi(r)|^2, \quad (4.28)$$

where the spatial distribution of the $c\bar{c}$ bound state is given by the squared wave function $|\phi_\psi(r)|^2$. The function $R(r)$ in eq. (4.28) describes the resolution capability of the colour field inside the light hadron. Its wave length is of order $\Lambda_{\text{QCD}}^{-1}$, and so it cannot resolve the charge content of very much smaller bound states; in other words, it does not “see” the heavy quarks in a bound state of radius $r_Q \ll \Lambda_{\text{QCD}}^{-1}$ and hence cannot rearrange bonds. The resolution $R(r)$ will approach unity for $r\Lambda_{\text{QCD}} \gg 1$ and drop very rapidly with r for $r\Lambda_{\text{QCD}} < 1$, in the functional form

$$R(r) \simeq (r\Lambda_{\text{QCD}})^n, \quad r\Lambda_{\text{QCD}} < 1, \quad (4.29)$$

with $n = 2$ [49] or 3 [39]. As a result, the integrand of eq. (4.28) will peak at some distance r_0 , with $r_Q < r_0 < \Lambda_{\text{QCD}}$. Since the bound state radius of the quarkonium ground state decreases with increasing heavy quark mass, while $R(r)$ is m_Q -independent, $r_0 \rightarrow r_Q \rightarrow 0$ as $m_Q \rightarrow \infty$. Hence P_r vanishes in the limit

$m_Q \rightarrow \infty$ because $R(r_0)$ does, indicating that the light quarks can no longer resolve the small heavy quark bound state.

Fig. 4.3: Rearrangement transition $\rho + J/\psi \rightarrow D + \bar{D}$

In a potential picture, the situation just described means that the charm quarks inside the J/ψ have to tunnel from $r = r_\psi$ out to a distance at which the light quarks can resolve them, i.e., out to some $r \simeq c\Lambda_{\text{QCD}}^{-1}$, where c is a constant of order unity (Fig. 4.4). Such tunneling processes are therefore truly non-perturbative: they cover a large space-time region, of linear size $\Lambda_{\text{QCD}}^{-1}$, and do not involve any hard interactions. Following [48], we shall here estimate the contribution of non-perturbative tunnelling to the dissociation of quarkonium states.

Fig. 4.4: J/ψ dissociation by tunneling

In general, the problem of quark tunneling cannot be solved in a rigorous way, since it involves genuine non-perturbative QCD dynamics. However, the large mass of the heavy quark allows a very important simplification, the use of the quasiclassical approximation. In this approximation, the rate of tunneling R_{tun} can be written down in a particularly transparent way: it is simply the product of the frequency $\omega_{\psi'}$ of the heavy quark motion in the potential well and the tunnelling probability P_{tun} when the quark hits the wall of the well,

$$R_{\text{tun}} = \omega_{\psi'} P_{\text{tun}} \quad (4.30)$$

The frequency ω_{ψ} is determined by the gap to the first radial excitation,

$$\omega_{\psi} \simeq (M_{\psi'} - M_{\psi}) \simeq E_{\psi}. \quad (4.31)$$

Consider now the potential seen by the $c\bar{c}$ (Fig. 4.4). For a particle of energy E , the probability of tunneling through the potential barrier $V(r)$ is obtained from the squared wave function in the “forbidden” region. It can be expressed in terms of the action W calculated along the quasiclassical trajectory,

$$P_{\text{tun}} = e^{-2W}, \quad (4.32)$$

where

$$W = \int_{r_1}^{r_2} |p| dr, \quad (4.33)$$

Here the momentum $|p|$ is given by

$$|p| = [2M(V(x) - E)]^{1/2}, \quad (4.34)$$

and r_1, r_2 are the turning points of the classical motion determined from the condition $V(r_i) = E$.

In our case, the width of the barrier is approximately $0.6 \Lambda_{\text{QCD}}^{-1}$, while its height $(V - E)$ is equal to the dissociation threshold E_Q . The mass M in Eq.(4.34) is the reduced mass, $M = m_Q/2$. We thus have

$$W \simeq 0.6 \sqrt{m_Q E_Q} / \Lambda_{\text{QCD}}. \quad (4.35)$$

For the J/ψ , we get from (4.35) the value $W \simeq 3$; this *a posteriori* justifies the use of quasiclassical approximation, which requires $S > 1$.

Using eq. (4.35), we obtain as final form for the tunneling rate (4.30)

$$R_{\text{tun}} = E_{\psi} \exp - (1.2 \sqrt{m_c E_{\psi}} / \Lambda_{\text{QCD}}). \quad (4.36)$$

With the above mentioned J/ψ parameters, this leads to the very small dissociation rate

$$R_{\text{tun}} \simeq 9.0 \times 10^{-3} \text{ fm}^{-1}. \quad (4.37)$$

In terms of R_0 , the J/ψ survival probability is given by

$$S_{\text{tun}} = \exp \left\{ - \int_0^{t_{\text{max}}} dt R_{\text{tun}} \right\}, \quad (4.38)$$

where t_{max} denotes the maximum time the J/ψ spends adjacent to the light hadron. In the limit $t_{\text{max}} \rightarrow \infty$, $S_{\psi'}$ vanishes. However, the uncertainty relations prevent a localisation of the two systems in the same spatial area for long times. From $\Delta x \leq \Lambda_{\text{QCD}}^{-1}$ we get $\Delta p \geq \Lambda_{\text{QCD}}$, so that the longest time which the J/ψ can spend in the interaction range of the light hadron is

$$t_{\text{max}} = \Lambda_{\text{QCD}}^{-1} \left(1 + \frac{m^2}{\Lambda_{\text{QCD}}^2} \right)^{1/2}, \quad (4.39)$$

with m for the mass of the light hadron. For nucleons or vector mesons, this time is 4 - 5 fm, and with this, the survival probability is very close to unity; hence non-perturbative tunneling interactions provide only negligible contributions to J/ψ dissociation.

In addition to such tunnelling, there can be direct and sequential thermal excitation to the continuum. In particular the latter still requires further analysis [48].

4.3 The Experimental Study of Charmonium-Hadron Interactions

We now return to the break-up rate for J/ψ interactions with slow light hadrons, as predicted by heavy quark QCD. Since this prediction is crucial for the use of quarkonia as confinement/deconfinement probe, it certainly must be checked experimentally. We shall now consider how that can be done.

In Eq. (3.22), we had obtained the survival probability for charmonium states i in $p-A$ interactions in terms of the break-up cross sections σ_i . We now just have to insert the cross section (4.24) into this expression to obtain the J/ψ survival probability as function of x_F . In Fig. 4.5 the result is shown for a cms collision energy of 17.4 GeV, the value provided in CERN SPS Pb -beam experiments. Included in this figure is also the colour octet interaction region $x_F \geq 0$, together with data [27] and the fit obtained in section 3.2. We see that in the regime $x_F \leq -0.2$, in which physical J/ψ 's interact with the nuclear medium, the survival probability is essentially unity, due to the smallness of the cross section (4.24) in the threshold region.

This picture has to be contrasted to the geometric absorption approach, which provides the basis for all hadronic accounts of J/ψ -suppression in nucleus-nucleus collisions. Here the J/ψ cross section is assumed to attain its high energy value $\sigma_{J/\psi N}(s = \infty) \simeq 2.5$ mb immediately at threshold. The survival probability for this case is also shown in Fig 4.5, and is seen to differ both qualitatively and quantitatively from that based on the QCD result (4.24). Through a measurement of J/ψ -production in $p - A$ collisions at negative x_F , technically attainable through data from a Pb -beam incident on a hydrogen or deuterium target at positive x_F , one can thus check if the threshold behaviour predicted by short distance QCD for inelastic J/ψ -hadron interactions is indeed correct.

The actual experimental results will differ from what is shown in Fig. 4.5 for two reasons. We have already noted in section 3.4 the modifications expected because of the EMC effect. A second reason is that the J/ψ peak observed in the measured dilepton spectrum is only to about 60 % due to directly produced $1S$ $c\bar{c}$ J/ψ resonances; the remaining 40 % are mainly due to χ_c production with the subsequent radiative decay $\chi_c \rightarrow J/\psi + \gamma$ [22,23].*

Concerning the EMC effect, we can either fold it into the predictions shown in Fig. 4.5, or it can be measured independently and then removed from the J/ψ -production data. Measuring Drell-Yan dilepton production in the same kinematic region provides $R_{A/p}(x)$ (see Eq. (3.23)) for quarks directly, without any final state modification. A measurement of open charm production leads to $R_{A/p}(x)$ in just the same superposition of quark-antiquark annihilation and gluon fusion as in charmonium production, but again without any final state effect. Separate measurements of Drell-Yan and/or open charm production would thus determine the EMC modification without additional final state effects. Such measurements can therefore be used to unfold the EMC modification of J/ψ data.

To avoid any χ_c admixture, it would of course be ideal to measure both J/ψ and χ_c -production directly in $p - A$ collisions [23]. This may be too difficult due to the abundance of photons in the case of heavy nuclei as targets. We shall therefore try to estimate the effect of the χ_c admixture in the two scenarios considered here by simply adding the corresponding predictions with the noted 60/40 weights. For the geometric approach with its asymptotic cross sections, this requires a calculation completely analogous to that for the J/ψ , but now with $\sigma_\chi \simeq 6$ mb and the correspondingly longer resonance formation time in the \bar{L} of Eq. (3.22).

The application of short distance QCD to calculate the inelastic χ_c -hadron cross section is as reliable as for the J/ψ , since the binding energy of the χ_c , $\epsilon_\chi \simeq 0.24$ GeV, is about equal to Λ_{QCD} . Nevertheless, such a calculation can give us

* We shall for simplicity ignore here a further small contribution ($\lesssim 5$ %) from ψ' decay).

some idea of the expected behaviour. The result is [25]

$$\sigma_{\chi N}(\bar{s}) \simeq 11.3 \text{ mb} \times \left(1 - \left[\frac{2M_\chi(m + \epsilon_\chi)}{(\bar{s} - M_\chi^2)} \right] \right)^{6.5}. \quad (4.40)$$

The asymptotic value is thus a factor two larger than the geometric estimate; this is a consequence of the fact that short distance QCD [40,41] leads to higher powers in the bound state radii than just r^2 . The behaviour of the $\chi_c N$ cross section (4.40) is shown in Fig. 4.2. Inserting Eq. (4.40) into (3.22) and adding J/ψ and χ_c contributions then gives us the short distance QCD survival probability for the measured J/ψ . It and the corresponding result from the geometric approach are shown in Fig. 4.6. As seen, for $x_F \leq -0.4$, the two approaches differ qualitatively in their functional form and quantitatively by more than 20%. An experimental test of the theory for the interaction of heavy quarkonia with light hadrons should therefore be possible.

As last point in this section, we comment briefly on the interaction of the ψ' with light hadrons. Since its binding energy is only about 50 MeV, it lies almost at the open charm threshold and can definitely not be treated by short distance QCD. Here it might therefore not be unreasonable to assume that it attains its high energy value rather soon after threshold. This value can only be estimated by geometric arguments, and these suggest around 10 mb [46,47].

5. Quarkonium Dissociation in Confined and Deconfined Media

Here we first show that at fixed temperature (or energy density), a deconfined medium contains much harder gluons than a confined medium. Tightly bound quarkonia probe gluon hardness: while they were found to remain essentially unaffected in a confined medium at $T \lesssim 0.5$ GeV, a QGP of such temperature is shown to be very effective in their dissociation. To relate this to the environment produced in nuclear collisions, the resulting charmonium survival is studied in the case of isentropic longitudinal expansion.

5.1 The Parton Structure of Confined vs. Deconfined Matter

The ultimate constituents of matter are evidently always quarks and gluons. What we want to know is if these quarks and gluons are confined to hadrons or not. Let us therefore assume that we are given a macroscopic volume of static strongly interacting matter and have to analyse its confinement status.

As prototype for matter in a confined state, we consider an ideal gas of pions. Their momentum distribution is thermal, i.e., for temperatures not too low it is given by $\exp(-E_\pi/T) \simeq \exp(-p_\pi/T)$. Hence the average momentum of a pion in this medium is $\langle p_\pi \rangle = 3 T$. 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 $x = p_g/p_\pi$.¹⁾ As a consequence, the average momentum of a gluon in confined matter is given by

$$\langle p_g \rangle_{\text{conf}} = \frac{1}{5} \langle p_\pi \rangle = \frac{3}{5} T. \quad (5.1)$$

Hence in a medium of temperature $T \simeq 0.2$ GeV, the average gluon momentum is around 0.1 GeV. In contrast, the distribution of gluons in a deconfined medium is directly thermal, i.e., $\exp(-p_g/T)$, so that

$$\langle p_g \rangle_{\text{deconf}} = 3T. \quad (5.2)$$

Hence the average momentum of a gluon in a deconfined medium is five times higher than in a confined medium²⁾; for $T = 0.2$ GeV, it becomes 0.6 GeV. An immediate consequence of deconfinement is thus a considerable hardening of the gluon momentum distribution. Although we have here presented the argument for massless pions as hadrons, it remains essentially unchanged for heavier mesons (ρ/ω) or nucleons, where one can use a non-relativistic thermal distribution for temperatures up to about 0.5 GeV. We thus have to find a way to detect such a hardening of the gluon distribution in deconfined matter.

The lowest charmonium state 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 infra-structure, 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, and this was the basis of the short distance QCD calculations presented in the previous chapter.

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, $\langle p_g \rangle_{\text{conf}} \simeq 0.1$ GeV, which is too soft to resolve the J/ψ as a $c\bar{c}$ bound state and much less than the binding energy ϵ_ψ , so that the J/ψ survives;
- if the matter is deconfined, $\langle p_g \rangle_{\text{deconf}} \simeq 0.6$ GeV, which (with some spread in the momentum distribution) is hard enough to resolve the J/ψ and to break the binding, so that the J/ψ will disappear.

¹⁾ For very small x , recent results from HERA indicate a steeper increase; however, this does not affect our considerations here.

²⁾ 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.

The latter part of our result is in accord with the mentioned prediction that the formation of a QGP should lead to J/ψ suppression [1,16]. There it was argued that in a QGP, colour screening would prevent any resonance binding between the perturbatively produced c and \bar{c} , allowing the heavy quarks to separate. At the hadronisation point of the medium, they would then be too far apart to bind to a J/ψ and would therefore form a D and a \bar{D} . Although the details of such a picture agreed well with the observed J/ψ suppression [50], it seemed possible to obtain a similar suppression by absorption in a purely hadronic medium [51], through collisions of the type

$$J/\psi + h \rightarrow D + \bar{D} + X. \quad (5.3)$$

Taking into account the partonic substructure of such hadronic break-up processes, we now see that this is in fact not possible for hadrons of reasonable thermal momentum. Our picture thus not only provides a dynamical basis for J/ψ suppression by colour screening, but it also indicates that in fact J/ψ suppression in dense matter will occur *if and only if* there is deconfinement. We note, however, that the dynamical approach to J/ψ suppression does not require a thermal equilibration of the interacting gluons, so that it will remain applicable even in deconfined pre-equilibrium stages.

While we have studied the hadronic part of the argument in detail in the previous sections, we have so far not considered the dynamics of quarkonium dissociation in a deconfined medium. This will be taken up in the next section.

5.2 Quarkonium Dissociation by Deconfined Gluons

In section 4.2 we had obtained the cross section for the dissociation of a tightly bound quarkonium by an incident light hadron. Eq. (4.23) was obtained [2] by convolution of the inelastic gluon-charmonium cross section with the gluon distribution in the light hadron. The gluon-quarkonium cross section itself is given by

$$\sigma_{g\Phi}(k) = \frac{2\pi}{3} \left(\frac{32}{3}\right)^2 \left(\frac{m_Q}{\epsilon_0}\right)^{1/2} \frac{1}{m_Q^2} \frac{(k/\epsilon_0 - 1)^{3/2}}{(k/\epsilon_0)^5}, \quad (5.4)$$

with k denoting the momentum of the gluon incident on a stationary quarkonium. In Fig. 5.1 we show the resulting break up cross section for gluon- J/ψ and gluon- Υ interactions as function of the gluon momentum k . They are seen to be broadly peaked in the range $0.7 \leq k \leq 1.7$ GeV for the J/ψ , with a maximum value of about 3 mb, and in the range $1.2 \leq k \leq 2.2$ GeV for the Υ , with a maximum of about 0.45 mb. Also shown in Fig. 5.1 are the corresponding cross sections for incident pions (note that now $k = 3$ in Eq. (4.2)), with high energy values of 3 mb and 0.5 mb for J/ψ and Υ , respectively. The break-up by incident hadrons is seen to be negligible up to momenta of around 4 GeV for the J/ψ and 7 GeV for the Υ . Fig.

5.1 thus provides the basis for the claim that in matter temperature $T \leq 0.5$ GeV, gluons of thermal momentum can break up charmonia, while hadrons cannot. We note here that, just as in the photoelectric dissociation of atoms, the break-up is most effective when the momentum of the gluon is somewhat above the binding energy. Gluons of lower momenta can neither resolve the constituents in the bound state nor raise them up to the continuum; on the other hand, those of much higher momenta do not see the (by their scales) large object and just pass through it.

To illustrate this more explicitly, we calculate the break-up cross section for the J/ψ as function of the temperature T of an ideal QGP. Using Eq. (5.4) with $m_c = 1.5$ GeV and the J/ψ binding energy of 0.64 GeV, we then get

$$\sigma_{gJ/\psi}(T) = 65 \text{ mb} \times \frac{\int_{\epsilon_0}^{\infty} dk k^2 e^{-k/T} (k/\epsilon_0 - 1)^{3/2} (k/\epsilon_0)^5}{\int_{\epsilon_0}^{\infty} dk k^2 e^{-k/T}}. \quad (5.5)$$

The result is shown in Fig. 5.2 and confirms that up to about $T \sim 0.5$ GeV, only a deconfined medium can dissociate J/ψ 's. We see moreover that the effective cross section for break-up in the temperature range $0.2 \leq T \leq 0.5$ GeV is about 1.2 mb. It is this value which will determine the suppression of the (pure 1S) J/ψ in a deconfined medium.

Before we turn to charmonium production in a more realistic non-static environment, we want to consider briefly the possible role of quarks in the dissociation of charmonia in a QGP. This can be treated in a fashion quite similar to the interaction of quarkonia with light hadrons; the gluon distribution function in the hadron should now be replaced by an effective gluon distribution “in” quarks, i.e., the quark splitting function $P(x)$ characterising the process $q \rightarrow q + g$. It's functional form for $x \rightarrow 1$ is fixed both by quark counting rules [52] and the Altarelli-Parisi equations [53],

$$P(x) \sim (1 - x)^2, \quad (5.6)$$

which leads to an average gluon momentum

$$\langle p_g \rangle_{\text{deconf}}^q = \frac{3}{4} T. \quad (5.7)$$

The average momentum of gluons emitted by thermal quarks is thus higher than for gluons confined to thermal hadrons; it is nevertheless low enough to suggest that in an equilibrium plasma, the direct interaction with thermal gluons is the main dissociation mechanism. In pre-equilibrium, however, the quarks can be much faster and become the dominant cause of dissociation.

5.3 Charmonium Survival in an Expanding Medium

The exponential quarkonium survival probability of the general form (3.20) applies to a stationary medium of finite size. In nuclear collisions, the medium

needs some proper time t_0 to be formed, and it then expands until a time t_f when the era of strong interactions in the medium ends. Hence slow quarkonia in the medium of not too high an initial temperature will in general stop interacting when the medium has cooled down enough, even though they have not yet left it. The survival probability of a quarkonium state i in such an expanding medium can be written as

$$S(T_0) = \exp\left\{-\int_{t_0}^{t_f} dt n(t) \sigma_i(t)\right\}, \quad (5.8)$$

where $n(t)$ is the density of scattering centers at time $t_0 \leq t \leq t_f$ and $\sigma_i(t)$ the break-up cross section for state i in the medium at that time. If we assume isentropic longitudinal expansion, density, temperature and time are related by

$$\frac{n(t_0)}{n(t)} = \left(\frac{T_0}{T}\right)^3 = \frac{t}{t_0}, \quad (5.9)$$

with $T \leq T_0$. Using this relation, we can rewrite Eq. (5.6) in the form

$$S(T_0) = \exp\left\{-3n_h(T_0)t_0 \int_{T_f}^{T_c} dT \frac{\sigma_i^h(T)}{T} - 3n_g(T_0)t_0 \int_{T_c}^{T_0} dT \frac{\sigma_i^g(T)}{T}\right\}, \quad (5.10)$$

where T_f is the temperature of the medium for which the break-up of state i stops. The first term in the exponent corresponds to dissociation by hadrons, the second by gluons. We had seen that the main contribution to the dissociation comes from gluon-quarkonium interactions, making $S(T_0) \simeq 1$ up to $T_0 \simeq T_c \simeq 0.15$ GeV. Hence by inserting the cross section (5.5) (Fig. (5.2)) into Eq. (5.10), we obtain the survival probability of a J/ψ in a medium in isentropic longitudinal expansion, for an initial temperature $T_0 \geq T_c$; below that, it is unity. The thermalisation time t_0 is generally argued to be around 1 fm; and for an ideal QGP, the density of gluons at temperature T is given by the Stefan-Boltzmann form

$$n(T) = \frac{8\pi^2}{45} T^3. \quad (5.11)$$

The resulting suppression as function of temperature is shown in Fig. 5.3. The survival probability is seen to decrease very rapidly with increasing temperature, essentially vanishing for $T \gtrsim 0.3$ GeV.

6. Pre-Equilibrium Deconfinement

In the previous chapter, we established quarkonium suppression as a probe to check if a given sample of strongly interacting *matter* consists of deconfined quarks and gluons. In an actual nuclear collision, however, such a suppression

could have been caused before the onset of thermalisation. Here we therefore study the possibility of pre-equilibrium quarkonium suppression.

6.1 Shadowing and J/ψ Suppression in Nucleus-Nucleus Collisions

Consider a nucleus-nucleus collision at CERN SPS energy ($\sqrt{s} \simeq 20$ GeV in the nucleon-nucleon cms), leading to J/ψ production at mid-rapidity, $y_{J/\psi}=0$, with the production mechanism as described in Chapter 2. During the first 0.25 fm, the produced $c\bar{c}$ pair is a colour octet and will interact as such with the passing nucleons of target and projectile. The momenta of these nucleons in the rest system of the J/ψ will be around 10 GeV. The situation seen by the J/ψ is thus very similar to that encountered at $x_F = 0$ in collisions of 200 GeV/c protons with a nuclear target. As discussed in Chapter 3, one here finds a suppression dominated by “nuclear shadowing”, i.e., destructive interference of the scattering on different nucleons. Such an effect will now arise from both projectile and target, however, and this suppression must be removed before final state effects can be studied [11].

To see the effect of this shadowing correction, we show in Fig. 6.1 the latest data taken by the NA38 collaboration at CERN in $p - U$ and $S - U$ collisions [54,55]. Here the suppression is measured with respect to the high mass Drell-Yan continuum, isospin corrected for $p - U$. The nucleus-nucleus data are shown as function of the neutral transverse hadronic energy E_T^0 produced in the collision. The $p-U$ value is simply included in this figure; it is not associated to any particular E_T^0 . The difference between the $p - p$ and $p - A$ results is, as discussed in Chapter 3, essentially given by the shadowing function $R_{A/p}^{sh}(x_F \simeq 0)$. To correct the data for shadowing effects, we therefore divide the $p - U$ data by $R_{U/p}^{sh}(0)$ and the $S - U$ data by $R_{S/p}^{sh}(0) \times R_{U/p}^{sh}(0)$. The result is shown in Fig. 6.2; the $p-U$ value and that for the S beam at low E_T^0 now agree. The remaining suppression in the nuclear collision data is now due to effects on the colour singlet $c\bar{c}$ in its state after the time of colour neutralisation.

We thus want to study the effect of the environment on this $c\bar{c}$. If we ignore nuclear stopping, target and projectile nucleons retain in the $c\bar{c}$ cms their initial momenta $P_N \simeq 10$ GeV/c. The $c\bar{c}$ then sees two nuclei, each Lorentz-contracted to about 1 fm, since $(2m/\sqrt{s}) \simeq 0.1$; for simplicity, we consider here $A - A$ collisions. At the colour neutralisation time $\tau_8 \simeq 0.25$ fm, the centers of these thin discs have become separated by approximately 0.5 fm, so that the target and projectile nucleons still have considerable overlap. Stopping would slow down the nucleons to increase this. The detailed kinematics encountered by the colour singlet $c\bar{c}$ is, however, quite unimportant for the essential question: can a pre-equilibrium state of hadrons account for the observed suppression? The initial medium, at the point of maximum overlap between target and projectile, is one of twice nuclear density, containing nucleons of 10 GeV/c momentum in the $c\bar{c}$ rest system; equivalently, we

can assume normal nuclear density and an effective $c\bar{c}$ path length $L = 2 \times 3R_A/4$, with $R_A = 1.12 A^{1/3}$ [34]. Whatever scattering processes now occur lead to an increase of the density n , but at the same time to a decrease of the momenta P of the hadronic constituents of the medium. Momentum conservation requires for the momentum flow through a given surface F

$$P n F = P_N n_0 F = \text{const.} \quad (6.1)$$

so that the momentum in fact drops as $1/n$. If we neglect the cross section reduction during the time needed for the singlet $c\bar{c}$ to become a full-sized resonance, the survival probability of a J/ψ in the medium becomes

$$S_{J/\psi} = \exp\{-n_0\sigma(P_N)L\}. \quad (6.2)$$

The medium actually seen by the J/ψ will have a higher density [56], but its constituents will have lower momenta. In view of Eq. (6.1), we can nevertheless use the form (6.2) to calculate the survival probability. In contrast to the phenomenological fit of [34], Eq. (6.2) contains the break-up cross section at the actual collision energy, and in Chapter 4 we have seen that this cross section is much below its geometric high energy value. At $P_N = 10$ GeV/c, we find (see Fig. 4.2) $\sigma_{J/\psi} \simeq 0.8$ mb. Inserting this into Eq. (6.2), with an average value $L \simeq 10$ fm for central $S - U$ collisions, leads to $S \simeq 0.9$ for the survival probability. Since we have here neglected the cross section reduction during the J/ψ evolution as well as the path length reduction by the part already included in the shadowing correction, the actual survival probability will be larger. Thus hadronic pre-equilibrium interactions cannot account for the measured J/ψ suppression.

We therefore conclude that the J/ψ suppression observed in the NA38 experiment [15] provides evidence for the existence of *deconfined* gluons in the medium probed by the J/ψ . By this we mean that the medium in which the J/ψ finds itself after a central nucleus-nucleus collision contain gluons whose momentum distribution is harder than that found in mesons or nucleons [57]. Only in collisions with such gluons can the tightly bound J/ψ be dissociated; gluons confined to hadrons are not sufficiently hard. We emphasize that this conclusion makes crucial use of the energy dependence of the J/ψ -light hadron break-up cross section as calculated in short distance QCD. As pointed out in Chapter 4, this result can and should be experimentally checked.

6.2 Deconfined Gluons and QGP

The existence of deconfined gluons is clearly *not* equivalent to the existence of a QGP, in which such gluons are in thermal equilibrium. It is only the first step towards a QGP: it shows that in the large-scale medium there exist deconfined

partons. Their thermalisation comes at the end of a parton interaction cascade, and it is not at all clear whether at present energies there are enough deconfined partons in a sufficiently large and long-lived system to reach this stage. How can we check experimentally whether parton thermalisation is or is not achieved?

The cascade formation of a QGP [58,59] starts with the production of gluons in primary collisions; these then interact again to produce secondary gluons, quite possibly through multigluon production [60], and so on, until production and absorption balance to form an equilibrium system. In equilibrium, the number of gluons per unit volume is simply determined by the energy (or entropy) density of the system. In the pre-equilibrium stage, it is lower and (still) proportional to the number of primary collisions. We thus expect J/ψ suppression by deconfined gluons in a pre-equilibrium system to increase with the number of nucleon-nucleon collisions (or equivalently, with the effective path length of the J/ψ). On the other hand, equilibrium suppression would be independent of the number of such collisions and depend on the effective energy density of the system only [57].

To test this, we can study $Pb - Pb$ collisions as function of increasing centrality (i.e., of increasing E_T). In this case, the energy density remains essentially unchanged [29] while the number of collisions increases.* If the J/ψ suppression is found to *increase*, the system cannot have reached equilibrium, and hence the suppression must be due to deconfined gluons in the pre-equilibrium stage. Once the suppression becomes independent of the number of primary collisions, equilibrium and hence the QGP is reached. The forthcoming NA50 Pb -beam data from the CERN-SPS should thus be able to resolve this question.

7. Conclusions

We first summarize the essential theoretical points of this work, and then survey the main conclusions we think can be drawn from present data.

Quarkonium production in hadron-hadron collisions is today quite well understood in terms of elementary parton interactions (gluon fusion, quark-antiquark annihilation). The distributions of the partons within the colliding hadrons are determined in deep inelastic lepton-hadron scattering, and the (non-perturbative) colour neutralisation of the produced heavy $Q\bar{Q}$ pair can also be fixed empirically.

Hadron-nucleus collisions determine what happens to quarkonium production in a confined medium. In the production of fast quarkonia in the nuclear rest frame, a fast colour octet $Q\bar{Q}$ passes the medium, leading to quantum-mechanical interference (“nuclear shadowing”) and energy loss. Slow quarkonia in the nuclear rest

* Since a change in centrality of $S - U$ collisions changes both the number of primary collisions and the effective energy density, this does not distinguish pre-equilibrium from equilibrium.

frame are subject to EMC effect modifications of the colliding partons in addition to collisions with nucleons in the nucleus.

Because of the small size and the large binding energy of the lowest quarkonium states, their interaction with light hadrons is calculable in short distance QCD. They can interact in leading order only through the exchange of a hard gluon, and the gluon distribution in the light hadrons is known to be very soft. The resulting prediction is a cross section which rises very slowly from threshold to its high energy value, suppressing strongly any break-up of quarkonium ground states by slow mesons or nucleons.

As a consequence of this suppressed quarkonium dissociation by light hadrons, confined matter (at temperatures $T \simeq 0.5$ GeV) becomes transparent to J/ψ 's. The momentum of deconfined thermal gluons, on the other hand, is large enough to give rise to effective J/ψ dissociation; such dissociation can occur also by deconfined gluons not in equilibrium. Strongly interacting matter thus leads to J/ψ suppression if and only if it is deconfined. The loosely bound ψ' can be broken up in both confined and deconfined matter, though presumably more in a deconfined medium.

In nucleus-nucleus collisions, the interaction of the coloured $c\bar{c}$ pair with target and projectile nucleons can be taken into account through the nuclear shadowing determined in hadron-nucleus collisions. Any strong J/ψ suppression remaining after removal of these shadowing effects is accountable only by interaction with deconfined gluons.

What have we then learned from the $h - A$ [27,28] and $A - B$ [15,61] data available so far? What should the forthcoming Pb -beam data clarify?

The equality of J/ψ and ψ' suppression in $h - A$ collisions for $x_F \geq 0$, as well as the size and x_F dependence of the observed effect, are in full accord with the passage of a colour octet through nuclear matter. The equality of J/ψ and ψ' suppression and the observed x_F dependence are in clear disagreement with any description based on the absorption of physical charmonium states in nuclear matter.

There is a lack of data for charmonium production in a kinematic regime in which fully formed J/ψ 's could interact with nuclear matter. Such data could be obtained by experiments using the Pb -beam incident on a light target [25].

The observed difference between J/ψ and ψ' suppression in nucleus-nucleus interactions [61] indicates that the relevant mechanism here is different from that in $h - A$ collisions.

The observed J/ψ suppression cannot be accounted for in terms of nuclear shadowing on both projectile and target; even after removal of shadowing, an E_T -dependent suppression (about 40 % between low and high E_T) remains.

If the strong threshold suppression of J/ψ break-up by light hadrons, as predicted in short distance QCD, is experimentally confirmed, the J/ψ suppression

observed in $O - U$ and $S - U$ interactions can only be accounted for by the presence of deconfined gluons.

It remains open whether the deconfined gluons required for J/ψ suppression in nuclear collisions are already equilibrated; hence their existence does not establish QGP formation, but quite likely only a first step towards a thermalised deconfined medium. If the gluons are not in equilibrium, the resulting J/ψ suppression should increase with increasing E_T in $Pb - Pb$ collisions at fixed energy; in equilibrium, the suppression would remain approximately constant.

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