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PROBING THE PRIMORDIAL STATE IN HIGH ENERGY HEAVY ION COLLISIONS*

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ABSTRACT

We first discuss the general problem of probing the early stages of dense matter produced in high energy heavy ion collisions, and then adress in particular the use of charmonium and bottonium spectra for this purpose. Removing initial state effects on the $c\bar{c}/b\bar{b}$ production process (gluon shadowing, initial state parton scattering), we obtain the relevant final state suppression form for the spectra. We then show how this can be used to test whether the observed J/ψ suppression in nuclear collisions is due to colour deconfinement or absorption in dense strongly interacting matter.

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1. Introduction

1.1. The Analysis of Dense Matter

The basic aim of high energy heavy ion experiments is the analysis of strongly interacting matter. We want to study bulk systems in a density region in which statistical QCD provides the relevant theoretical description. From QCD we know that in sufficiently dense matter, colour screening will lead to quark deconfinement, and as a result, there will be a transition from hadronic matter to a plasma of deconfined quarks and gluons. The study of this predicted new state of matter, the quark-gluon plasma, is the ultimate aim of our work. Lattice QCD calculations [1] indicate that the transition will occur at a temperature $T_c \simeq 150$ MeV; this means that energy densities $\epsilon \gtrsim \epsilon_c \simeq 1$ GeV/fm^3 are necessary to achieve deconfinement. Cosmological estimates put the density of our universe above this value for the first 20 μs after the big bang. We are thus trying to recreate locally the conditions of the very early universe.

This endeavour leads to four crucial questions. We hope that the collision of two heavy nuclei leads to the production of droplets of energy deposited in the vacuum; they are the candidates for our mini-universe. We thus have to ask:

(1) Were these droplets *hot* enough? The initial energy density ϵ_o can be estimated from the emitted hadron energy per unit rapidity. For central collisions at the SPS, this already gives values of 2 GeV/fm^3 or more, and at RHIC or LHC we should reach energy densities of five to ten times ϵ_c [2].

(2) Were the droplets *large* enough? Since we want to observe critical phenomena (deconfinement and chiral symmetry restoration), we obviously need large systems. Hadron interferometry provides us with information about the system size at freeze-out, when strong interactions stop [3]. For light ion beams at the SPS (^{32}S), this gives volumes of around 500 fm^3. Extrapolations to Pb-Pb collisions at the SPS suggest $V_f \gtrsim 7000$ fm^3, and for RHIC and LHC we get $V_f \simeq 10^4 - 10^5$ fm^3. We thus expect that also in the regime of $\epsilon \gtrsim \epsilon_c$ the size of the droplets is very much larger than a typical hadronic volume of 1 - 2 fm^3.

(3) Was the system *thermal*? Even a large, dense system could be just a superposition of nucleon-nucleon collisions, far away from any thermal equilibrium state. To attain equilibrium, we need sufficient multiple scattering and rescattering among secondaries as well as primaries. This can be tested by studying particle production ratios [4], and there are definite indications for a considerable amount of rescattering and hence a possible onset of thermalisation.

If the three points mentioned so far can all be answered affirmatively, we are left with the perhaps most difficult question:

(4) How can we test if there was a *deconfined primordial state*? The hot dense matter produced in the collision immediately expands and hence cools off; eventually it hadronizes, and when the resulting hadrons are measured, the plasma is long gone. If the evolution of the primordial droplets follows equilibrium thermodynamics, all memory of the previous history is lost in this evolution, and by measuring the secondary hadrons, we cannot learn anything about a possible prior plasma phase.

We thus have to look for specific *non-hadronic* primordial state probes. Two kinds of such probes have so far been studied in detail.

One possibility is to study signals which were emitted in the primordial state and are not affected by the subsequent evolution of the system. The most promising candidates for such signals are thermal dileptons and photons [5]. However, they are more suitable as thermometers to measure the primordial temperature than as probes for the nature of the primordial state: the thermal spectrum of real or virtual photons is functionally the same for emission from a hadron gas and from a quark-gluon plasma.

Another approach is to consider probes which were present *before* any dense matter was formed and which are sensitive to deconfinement; their subsequent behaviour should then tell us something about the nature of the primordial state. The first such signal to be studied was J/ψ production in nuclear collisions [6]. The production of hard, perturbative jets could serve as a similar probe, but so far no really clear-cut predictions have been obtained [7]. The production of J/ψ and ψ' is readily accessible to present experiments, and it does show clear nuclear effects. In addition, at RHIC and LHC energies, a comparative study of charmonium and bottonium production should become possible. We shall therefore concentrate on the interpretation of this signal as a probe for the primordial state in nuclear collisions.

1.2. Charmonium and Bottonium Spectroscopy in a Screening Medium

The potential binding $c\bar{c}$ or $b\bar{b}$ systems to charmonium or bottonium state is in the absence of a medium given by [8]

$$V(r) \;=\; \sigma r - \left(\frac{\alpha}{r}\right), \tag{1}$$

where σ denotes the string tension and r the separation between quark and antiquark; the term (α/r) describes the Coulomb part of the potential as well as transverse string vibrations. In a medium, the presence of other colour charges will lead to a screening of the colour force; the potential now becomes [9]

$$V(r,\mu) \;=\; \sigma r \left[\frac{1-e^{-\mu r}}{\mu r}\right] \;-\; \left(\frac{\alpha}{r}\right) e^{-\mu r}; \tag{2}$$

here μ is the screening mass (the inverse "Debye screening length" for the colour force). Solving the Schroedinger equation

$$\left[2M - \left(\frac{1}{M}\right)\nabla + V(r,\mu)\right]\Phi(r,\mu) \;=\; E(\mu)\,\Phi(r,\mu) \tag{3}$$

with charm or bottom quark mass M then gives us for $\mu = 0$ the masses and radii of the different possible bound states, i.e., of the various charmonium and bottonium states. For sufficiently large μ, there are no more bound states, and we can determine the specific values of the screening mass, for which the different bound states dissolve. Lattice QCD provides a relation between screening mass

and energy density of the medium, so that we thus obtain the dissociation energy density ϵ_d for each resonant state. The result is shown in Table 1.

We see that all ϵ_d should be accessible already at the SPS, with the exception of that for the Υ. Conservative estimates give energy densities around 10 GeV/fm^3 for the LHC, so that the pure $1S$ $b\bar{b}$ ground state should remain bound even then.

We thus find that charmonium and bottonium states will melt in a deconfining medium. Can we observe this in nuclear collisions in the form of suppressed charmonium/bottonium production? In the next three sections, we will address different aspects of this question. In section 2, we compare deconfinement and absorption as different possible suppression mechanisms in dense matter. In section 3, we check to what extent a suppression can arise already from initial state effects, i.e., before any dense matter is formed, and in section 4 we summarize what remains of the experimentally observed J/ψ suppression once the initial state effects are removed. In closing, we discuss how one may use the resulting final state suppression data to probe the nature of the primordial state.

States :	ψ	ψ'	χ_c	Υ	Υ'	χ_b
M_x	3.1	3.7	3.5	9.6	10.0	9.9
r_x	0.22	0.44	0.35	0.12	0.25	0.20
ϵ_x	1.9	1.0	1.0	47.1	1.6	1.0

Table 1
Charmonium and bottonium masses M_x [GeV], radii r_x [fm], and dissociation energy densities ϵ_x [GeV/fm^3], from [10].

2. J/ψ Suppression in Dense Matter

The dominant mechanism for J/ψ production in hadronic collisions is gluon fusion* (Fig. 1): two gluons interact in a *hard* process (i.e,, described in terms of perturbative QCD) to form a $c\bar{c}$ pair; this radiates a gluon and subsequently undergoes *soft* (i.e., non-perturbative) binding to form the J/ψ, whose decay into a lepton pair can be observed in the measured dilepton spectrum. What effect can the presence of a dense medium have on this process?

In a deconfining medium, colour screening prevents the binding of the $c\bar{c}$ to form a J/ψ; hence the c and the $\bar{c}$ will simply "fly apart". At the time of hadronisation,

* We restrict ourselves here to production at small or intermediate x_F; for production at large x_F, other mechanisms come into play (see e.g. [11]).

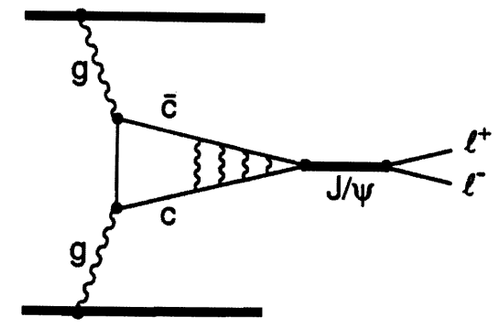

Fig. 1 J/ψ production via gluon fusion.

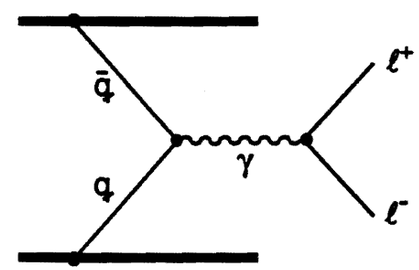

Fig. 2 Drell-Yan dilepton production via $q\bar{q}$ annihilation.

the two will thus be generally too far separated to bind with each other. All other hadrons share the same fate, of course; but in the resulting quark-gluon plasma, there are u, d and s quarks in sufficient numbers to lead via recombination to the production of the usual hadrons. The thermal production of further $c\bar{c}$ pairs in the medium is almost completely suppressed; at $T \simeq 200$ MeV, $\exp(-2m_c/T) \simeq 10^{-7}$. The c and the $\bar{c}$ from the early, hard process remain the only ones present, and if they are too far apart to couple with each other, they have to couple with the usual hadrons present and make a D and a $\bar{D}$. The presence of a deconfining medium thus deviates the J/ψ production channel into $D/\bar{D}$ production.

The Drell-Yan dilepton production rate (Fig. 2), on the other hand, is not affected by the medium, since the photons formed in $q\bar{q}$ annihilation can only interact electromagnetically with their surroundings. We thus conclude that the production of a colour deconfining medium in nuclear collisions will induce a suppression of the J/ψ peak in the dilepton spectrum, relative to the Drell-Yan continuum [6].

Following this prediction, the NA38 collaboration at CERN, in a series of O-Cu, O-U and S-U experiments [12], observed a J/ψ suppression of about 50% at the highest energy density (Fig. 3); the suppression was found to occur mainly for J/ψ's of low transverse momentum (Fig. 4). This triggered an intensive search for possible alternative suppression mechanisms (see [13,14] for surveys), which led to the conclusion that the observed behaviour could be accounted for as well by absorption in very dense hadronic matter. Let us therefore also consider this mechanism in some detail.

Deconfinement is a *global* phenomenon - the medium as a whole provides the screening of the colour force between the c and the $\bar{c}$. In contrast to this, absorption constitutes a *local* effect: by interactions of the form

$$J/\psi + x \rightarrow D + \bar{D} + \ldots \ , \tag{4}$$

the J/ψ is broken up when it collides with a single constituent x in the medium. Traversing a static medium of linear size L, the J/ψ has a survival probability $S = \exp{-L/\lambda}$, where λ is the mean free path of the J/ψ in the medium. It is given by $\lambda = 1/n\sigma$, with n denoting the density of constituents and σ the break-up cross-section.

The form (1) is evidently an oversimplification, since the actual medium is rapidly expanding and cooling. For longitudinal isentropic expansion, the survival probability becomes

$$S \ = \ \exp[-n_o t_o \int_{t_o}^{t_f} dt \ \sigma(t)/t \ ] \ \simeq \ \exp[-n_o t_o \sigma_{J/\psi} \ln(t_f/t_o) \ ]. \tag{5}$$

Here t_o denotes the thermalisation time for the medium; we take it to be about one fermi and hence roughly equal to the J/ψ formation time; we therefore assume $\sigma(t) \simeq \sigma_{J/\psi}$ is the J/ψ break-up cross-section. Further, n_o is the constituent density at $t = t_o$, and t_f denotes the freeze-out time, at which the medium has become too dilute for further interactions.

To compare this form with data, we have to identify the constituents. They are generally taken to be hadrons, and then the resulting break-up cross-section is

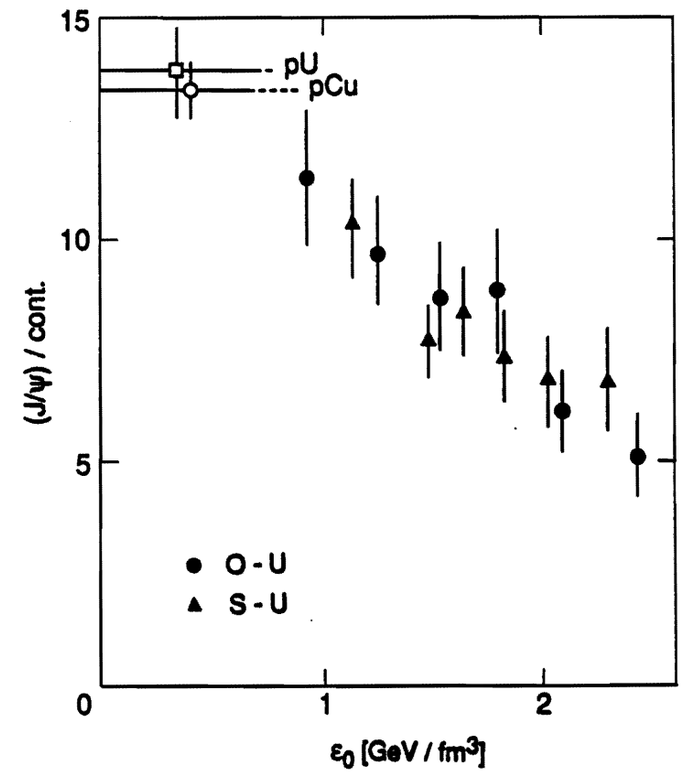

Fig. 3 J/ψ **suppression in nucleus-nucleus collisions as a function of the energy density, from [12,28].**

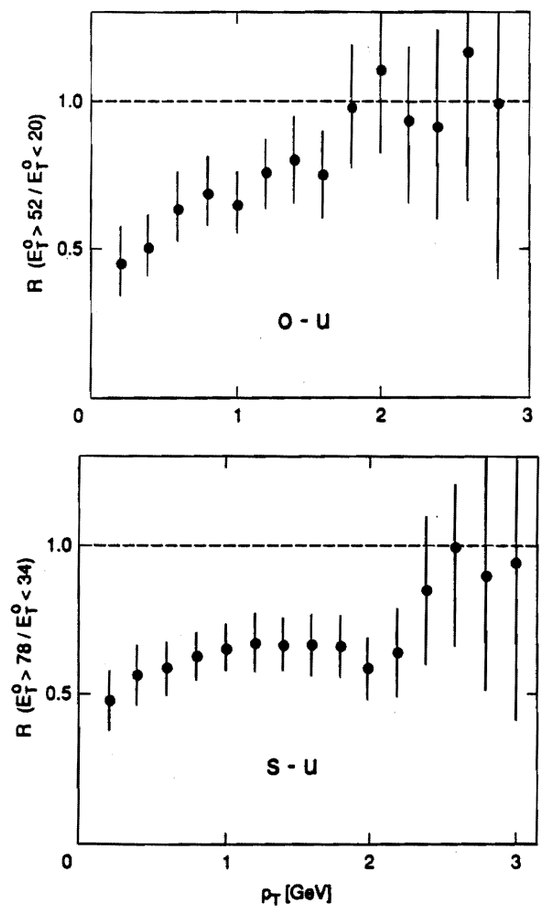

Fig. 4 J/ψ suppression in nucleus-nucleus collisions as function of the transverse momentum, from [12,28]. The function R shows the ratio of J/ψ production at high E_T to J/ψ production at low E_T, normalized to the same ratio for the Drell-Yan continuum.

about 2-4 mb, obtained by using geometric arguments [15]. The initial density is not so easy to estimate, since many of the observed secondaries are decay products of hadronic resonances. The highest initial energy density in the NA38 experiments lies around 2-3 GeV/fm^3; this would correspond to about 4-6 typical pions/fm^3. By taking resonance production into account, this can be reduced considerably, and one finds that $n_o \gtrsim 2$ hadrons/fm^3 produces a suppression of the observed magnitude. There is clearly an interplay between the choice of the size of the absorbing volume, the break-up cross-section, and the initial density. Reasonable estimates result in the mentioned value of n_o [16] – and even this still has about four hadrons squeezed into one typical hadronic volume (2 fm^3). The conceptual basis for such a picture remains unclear; nevertheless, we can conclude that absorption in a super-dense medium does account for the observed amount of J/ψ suppression [17,18,19]. We thus have two competing scenarios for the fate of the J/ψ in dense matter.

One obvious task for further studies is clearly the search for characteristics which distinguish the suppression by colour deconfinement from that by absorption. But before addressing this issue, we have to ask whether there are additional suppression effects in nuclear collisions due to phenomena taking place *before* the formation of dense matter.

3. Initial State Effects*

The description of heavy-flavour resonance production in QCD involves three sets of inputs— the distribution of the gluons in the initial hadrons or nuclei, the cross section for these to produce heavy-flavour quarks, and a model for the binding of the produced quarks and antiquarks to charmonium or bottonium states (see Fig. 1). Denoting these three elements symbolically by G, $d\hat{\sigma}$ and H, respectively, the hadroproduction cross section is given by the convolution

$$d\sigma(s, M, x_F) = \left(\sum_{ij} G_i(x_1) G_j(x_2) d\hat{\sigma}_{ij} \right) H(P)\, d\Gamma(P). \tag{6}$$

Here the variables x_1 and x_2 denote the fractional momenta carried by partons from the beam and the target, respectively. The heavy flavour pair has 4-momentum P, which we take also to be the momentum of the observed state, since the resonance is produced by a soft process. The Lorentz invariant phase space volume is denoted by $d\Gamma$; $\sqrt{s}$ is the CMS energy of the colliding hadrons, M the mass of the resonance and x_F its scaled longitudinal momentum. The different variables are related through the kinematical identities

$$x_1 = \frac{1}{2}\left(\sqrt{4\tau + x_F^2} + x_F\right), \quad x_2 = \frac{1}{2}\left(\sqrt{4\tau + x_F^2} - x_F\right), \quad \tau = M^2/s. \tag{7}$$

where we have for simplicity assumed $P_T \simeq 0$. Before we can study the effect of any dense medium on H, we first have to take into account any nuclear modifications of

* This and the following section are based largely on joint work with Sourendu Gupta, published in [20,21].

the structure functions $F(x)$. Such modifications are in fact expected, since they are observed in both J/ψ and Drell-Yan production induced by p-A interactions, where dense matter doesn't enter. In particular, the presence of very soft partons ($x \leq 0.1$) is suppressed in nuclei in comparison to nucleons (nuclear shadowing, see Fig. 5); it is also suppressed in the intermediate x-range ($0.2 \leq x \leq 0.6$) (EMC effect), and the P_T distribution of heavy resonances in collisions with nuclei are broader than in p-p collisions (initial state parton scattering). If any of these effects becomes important in nucleus-nucleus collisions, they have to be taken into account before we can use charmomium or bottonium production to probe primordial matter.

Let us begin with gluon shadowing [20]. It is found that J/ψ production (integrated over P_T) in p-A collisions, in comparison to the production in p-p interactions, is suppressed with increasing A; the same is true for ψ' and Υ production. This is seen in Fig. 6, where the ratio $R^i_{A/p} \equiv \sigma^i_{pA}/A\sigma^i_{pp}$ is shown for the different resonances i [22,23]. If the suppression is purely a structure function effect, we have

$$R_{A/p}(x_F, \sqrt{s}, M_i) = G_A(x_2)/G_p(x_2). \tag{8}$$

Hence the ratio, which *a priori* could depend on the three variables indicated, becomes a function of the one dimensionless variable $x_2 = (1/2)(\sqrt{4\tau + x_F^2} - x_F)$ only, with $\tau = M_i^2/s$, for all different states i (we here neglect possible small scale-breaking terms). Such a form can of course be valid only in the kinematic region in which the factorized gluon fusion form (6) is correct; this means that we have to exclude data at large x_F, where, as already noted, other mechanisms come into play [11]. We therefore restrict ourselves to the range $x_F \leq 0.6$ and show in Fig. 7 the data of the NA3 collaboration at CERN [22] for J/ψ production on p and Pt targets with 200 and 280 GeV p-beams, together with the data for J/ψ, ψ', and Υ production from the Fermilab collaboration E772 [23], on deuterium and W targets with a p-beam of 800 GeV. We see that all data fall quite well on one scaling curve and therefore conclude that the bulk of the suppression seen in Fig. 5 can be understood in terms of gluon shadowing. Data on other targets [23] and from deep inelastic muon scattering [24] are in accord with this.

In the region $x_2 \leq 0.1$, we can parametrize the data shown in Fig. 7 by the form

$$R_{A/p} = 1 + a \ln x_2 \ln A, \tag{9}$$

getting $a = 0.021 \pm 0.001$. If we take the form (9) to be the first term in the expansion of a power-law behaviour, we get

$$R_{A/p} \simeq A^{\alpha_{eff}-1}, \tag{10}$$

but with a power $\alpha_{eff} = 1 + a \ln[(1/2)(\sqrt{4\tau + x_F^2} - x_F)]$ which depends on x_F as well as on $\sqrt{s}$. This power agrees very well with the fitted values obtained at the different incident beam energies and for the different charmonium and bottonium states observed [22,23]. In section 4, we shall use the parametrization (9) to remove the effect of gluon shadowing from the measured nucleus-nucleus data.

We now turn to the effect of initial state parton scattering. In nuclear collisions, the gluons which eventually fuse to produce a $c\bar{c}$ pair can in principle have interacted

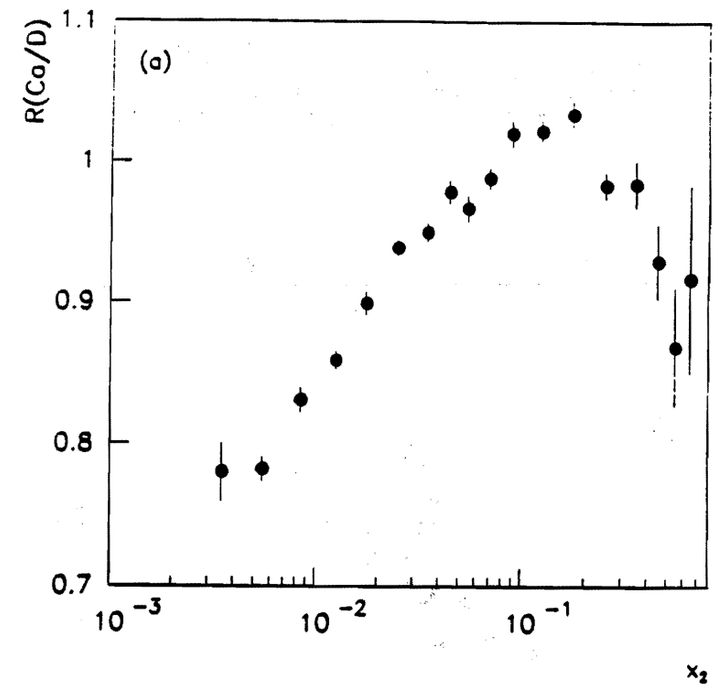

Fig. 5 Quark shadowing in deep inelastic muon scattering on calcium and deuterium targets [24].

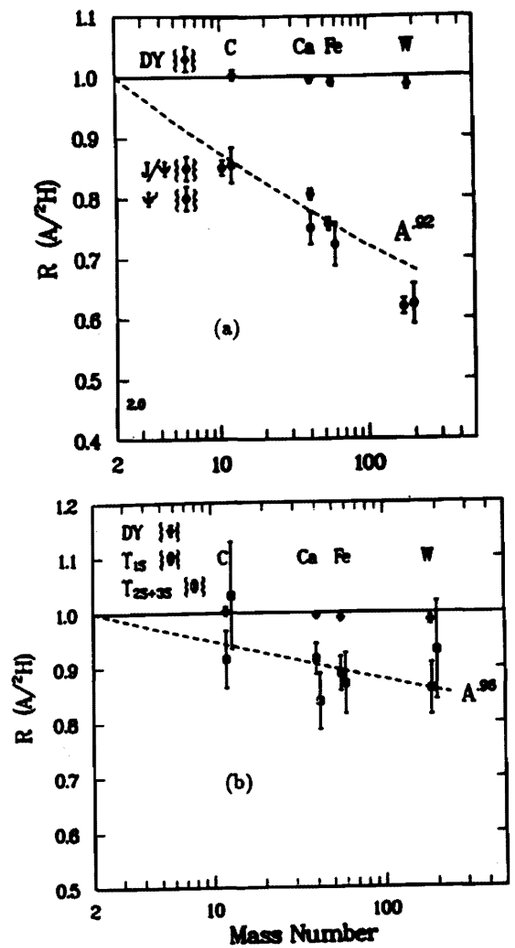

Fig. 6 (a) J/ψ and ψ' production in p-A collisions compared to that in p-p collisions, from [23]; (b) the same, for Υ production.

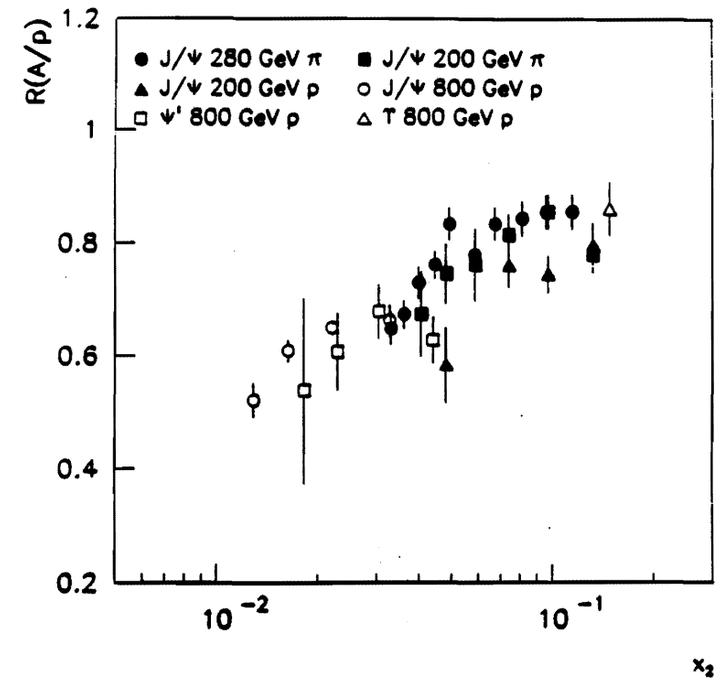

Fig. 7 Pion (200 and 280 GeV) and proton (200 and 800 GeV) data for the ratio of J/ψ, ψ' and Υ production on platinum [22] or tungsten [23] targets to that on hydrogen [22] or deuterium [23].

with other partons before this fusion. This will rotate the axis of the parton-parton fusion system relative to the beam axis and hence lead to a broadening of the P_T distribution for J/ψ and Drell-Yan dilepton production in p-A collisions, compared to that in p-p interactions [25–27].

Consider the result of this on the mean squared transverse momentum of J/ψ's,

$$\Delta_{\rm pA} \equiv \langle P_T^2 \rangle_{\rm pA} - \langle P_T^2 \rangle_{\rm pp}. \tag{11}$$

We expect $\Delta_{\rm pA} = \langle q^2 \rangle L_{\rm A}/\lambda$, where $\langle q^2 \rangle$ is the mean squared transverse momentum a parton acquires per collision, $L_{\rm A}$ the "pre-fusion" path length of the parton in the nucleus A, and λ its mean free path in A. Using $L_{\rm A} \sim R_{\rm A}$, with $R_{\rm A} = 1.2A^{1/3}$, we obtain $\Delta_{\rm pA} \simeq 0.048\ R_{\rm A}$ from fits to the J/ψ production in p-A collisions. With this and the assumption of a Gaussian distribution in the partonic transverse momentum q, we get

$$R_{\rm A/p}(x, P_T) = \left[\frac{G_{\rm A}(x_2)}{G_{\rm p}(x_2)}\right] \left[\frac{\langle P_T^2 \rangle_{\rm pp}}{\langle P_T^2 \rangle_{\rm pA}} \exp\left(P_T^2 \left[\frac{1}{\langle P_T^2 \rangle_{\rm pp}} - \frac{1}{\langle P_T^2 \rangle_{\rm pA}}\right]\right)\right] \tag{12}$$

for the ratio of the double differential p-A and p-p cross-sections. It suppresses production at small P_T, and this, as we shall see, accounts for much of the P_T dependence observed in nucleus-nucleus collisions [25–27]. Again we have to ask for the behaviour that remains when this initial state effect is removed.

4. Final State J/ψ Suppression

The NA38 experiment at CERN provides the ratio of the J/ψ signal to the Drell-Yan continuum for p-Cu, p-U, O-Cu, O-U and S-U collisions [12,28]. Using the parametrization (9), we calculate the shadowing effects relative to p-p collisions in each of these cases. From

$$R_{\rm (AB)/(pp)} = \frac{g_{\rm A}(x_1) g_{\rm B}(x_2)}{g_{\rm p}(x_1) g_{\rm p}(x_2)} = R_{\rm A/p}(x_1) R_{\rm B/p}(x_2) \tag{13}$$

and Eq.(9), we get for $\sqrt{s} = 20$ GeV the results shown in Table 2. We have assumed $\langle x_F \rangle \simeq 0$, and thus $x_1 = x_2 = M/\sqrt{s}$. At this value of x_1 and x_2, the quark structure function ratios entering the Drell-Yan rates are approximately unity. Hence if we divide the data for the various channels by the corresponding values of Table 2, we obtain the suppression beyond that due to shadowing effects in the initial state.

In Fig. 8, the resulting final state suppression is shown as a function of the energy density, using the data given in [28]. In particular, we note:

- for the highest energy densities in nuclear collisions, there still is a J/ψ suppression of around 50 % relative to p-A and p-p collisions;
- at the lowest energy densities measured in nuclear collisions, i.e, around 1 GeV/fm^3, there is no suppression relative to p-A and p-p collisions, once shadowing is taken into account.

The pattern seen in Fig. 8 is thus very suggestive of a threshold for suppression around $\epsilon \simeq 1$ GeV/fm^3, as expected if quark deconfinement is the relevant mechanism.

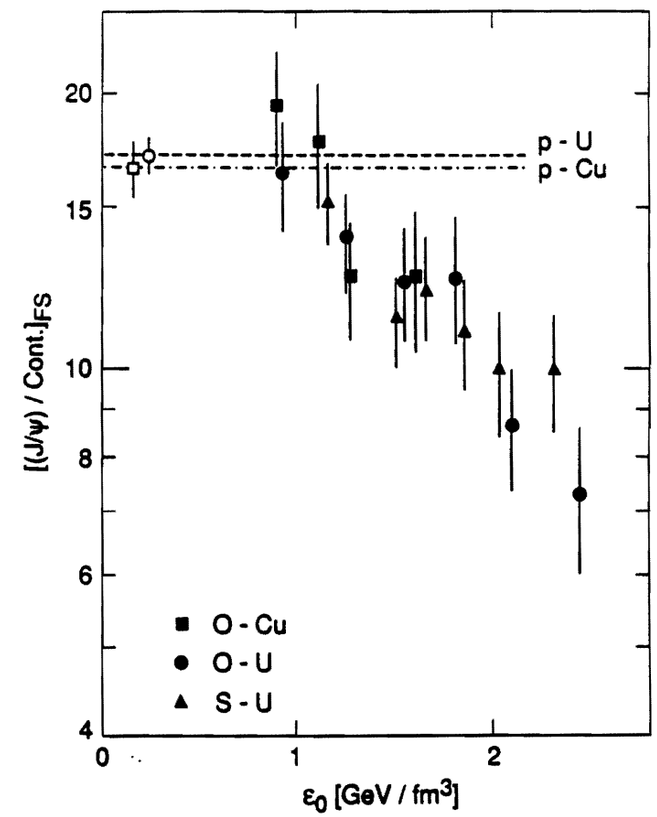

Fig. 8 Final state J/ψ suppression in nucleus-nucleus collisions as a function of the energy density.

Reaction	$R_{(A-B)/(p-U)}$
p − Cu	0.837
p − U	0.786
O − Cu	0.747
O − U	0.701
S − U	0.679

Table 2

Suppression by gluon shadowing in nuclear collisions at $\sqrt{s} = 20$ GeV, for $\langle x_f \rangle \simeq 0$, from [21].

We now turn to initial state parton scattering and its effect on the transverse momentum behaviour of the suppression [25–27]. In nucleus-nucleus collisions, the broadening arises due to the motion of the projectile through the target, as well as the other way around. Hence, using the same arguments as those following eq.(11), we get for central collisions of two large nuclei A and B

$$\Delta_{AB} \;=\; \langle P_T^2 \rangle_{AB} - \langle P_T^2 \rangle_{pp} \;\simeq\; 0.048\,(R_A + R_B). \tag{14}$$

The effective path length L_{AB} depends on the impact parameter b. It has its maximum value at $b = 0$ and decreases with increasing b, vanishing at $b = R_A + R_B$. This b-dependence translates into a dependence on the transverse hadronic energy E_T [25–27][29], since more central collisions lead to higher multiplicities and hence a higher E_T in the final state. The E_T dependence of Δ_{AB} for charmonium spectra from nucleus-nucleus collisions was shown to be fully accounted for by this effect [29]. We thus have to check what final state pattern for the tranverse momentum distribution itself remains when the "beam rotation" due to initial state transverse momentum broadening is included.

From NA38 we have the ratio $R_{H/L}$ of the Q_T distributions at high and at low E_T. To remove the effect of initial state parton scattering, $R_{H/L}$ has to be multiplied by

$$G_L/G_H \;\equiv\; \{\langle P_T^2 \rangle_H / \langle P_T^2 \rangle_L\} \exp\{-P_T^2 \left(\frac{1}{\langle P_T^2 \rangle_L} - \frac{1}{\langle P_T^2 \rangle_H} \right)\}. \tag{15}$$

The parameters $\langle P_T^2 \rangle_H$ and $\langle P_T^2 \rangle_L$ entering here can either be determined from geometric considerations relating E_T to impact parameter and path length, or they can be taken from the data. Since the data fit the geometric picture [29], we use them directly. In Fig. 9 we show the result for S-U collisions. As expected, the transverse momentum dependence of the suppression is now essentially removed; we have a constant suppression of around 60% up to $P_T \simeq 2-3$ GeV. For transverse momenta beyond this range, neither the factorized form (1) nor the Gaussian distribution in Eq.(8) makes any sense.

Let us summarize: we find that the final state J/ψ suppression in nuclear collisions, obtained after removing the effects of both gluon shadowing and initial state parton scattering, appears to set in rather abruptly at an energy density around 1 GeV/fm^3, increases to about 50% at the highest value of ϵ_o, and shows little dependence on the transverse momentum of the J/ψ. What consequences does this behaviour have for the attempts to account for the observed suppression in terms of either quark deconfinement by colour screening or absorption in dense hadronic matter?

Deconfinement sets in at a critical threshold energy density ϵ_c, although the threshold can be softened by quantum mechanical modifications of J/ψ formation [30]. Since sufficiently fast J/ψ's can escape from the deconfining medium in space or in time, deconfinement can also lead to a transverse momentum dependence [31]. The overall suppression observed by NA38 disappears for $P_T \simeq 2-3$ GeV; if interpreted in terms of an "escape" from the plasma, this leads to extremely short lifetimes for the deconfining medium [32,33]. We have seen, however, that once initial state parton scattering is removed, very little P_T dependence remains. This in turn removes the argument for very short plasma lifetimes. The surviving J/ψ's are in fact not those from $c\bar{c}$ pairs which are fast enough to escape from the medium before they have separated more than the diameter of the physical J/ψ, but rather those which are produced in the cooler edges of the interaction region. In addition, the mixed origin of the observed J/ψ has to be taken into account [34]: about 70% are produced directly as $1S$ $c\bar{c}$ states (we denote these states by ψ in what follows), the remainder come from χ_c decay [35]. While the χ_c melts essentially at the deconfinement point ϵ_c, the energy density needed for the melting of the ψ is about twice this value, as seen in Table 1. Hence much of the observed J/ψ suppression would come from a melting of the χ_c in the matter produced in present experiments. Let us see what this means quantitatively.

In the plasma scenario, the survival probability for the J/ψ depends only on the fraction of the interaction region "cool" enough to prevent melting. In a central collision of a small on a large nucleus one finds for nuclei of constant density an energy density profile of the form [10]

$$\epsilon_o(r) \;\sim\; \epsilon_o \left(1 - (r/R_B)^2\right)^{2/3}, \tag{16}$$

where R_B is the radius of the (smaller) projectile nucleus and r the radial distance from its centre. It leads to a survival probability

$$S_x(\epsilon_o) \;=\; [\Theta(\epsilon_x - \epsilon_o) + \Theta(\epsilon_o - \epsilon_x)(\epsilon_x/\epsilon_o)^{9/4}], \tag{17}$$

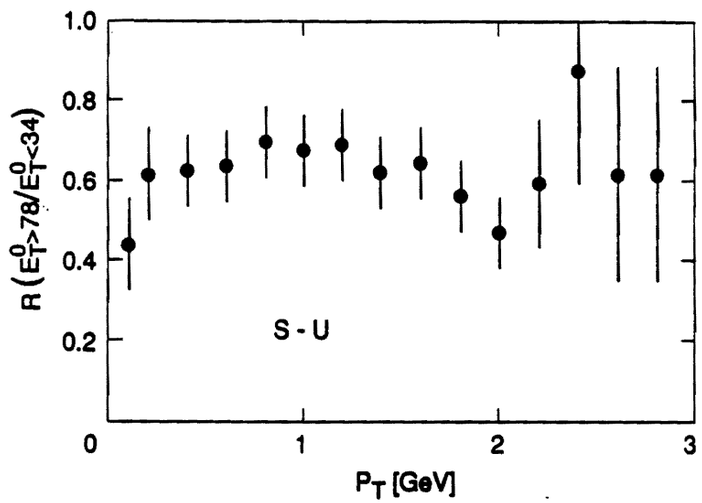

Fig. 9 Final state J/ψ suppression in $S-U$ collisions as a function of the transverse momentum.

where ϵ_x is the energy density needed to melt state x; with $\epsilon_\chi \simeq \epsilon_c$ and $\epsilon_\psi \simeq 1.9\ \epsilon_c$ from Table 1, we obtain

$$S_{J/\psi}(\epsilon_o) = 0.7\ [\Theta(1.9\ \epsilon_c - \epsilon_o) + \Theta(\epsilon_o - 1.9\ \epsilon_c)(1.9\ \epsilon_c/\epsilon_o)^{9/4}] + 0.3\ [\Theta(\epsilon_c - \epsilon_o) + \Theta(\epsilon_o - \epsilon_c)(\epsilon_c/\epsilon_o)^{9/4}] \tag{18}$$

as the (parameter-free) form of the survival function for the J/ψ. We see in Fig. 10 that it agrees fairly well with the present data. Note that if more than 30% of the observed J/ψ's come from χ_c decay, the survival probability will drop faster. It will, however, retain its bump structure around 2 GeV/fm^3, as long as the measured J/ψ comes from two channels with different dissociation temperatures.

In the range accessible to present experiments, absorption is quite independent of the transverse momentum of the J/ψ, and as we saw above, it can also give a sufficiently strong suppression. It cannot, however, lead to the threshold behaviour which seems to appear when gluon shadowing is removed (see Fig. 8); instead, we obtain a survival probability of the form (5). Using geometric estimates for the total cross sections [15], we have $\sigma_t(\psi) \simeq 3$ mb and $\sigma_t(\chi_c) \simeq 7$ mb. Taking into account the χ_c/ψ mixture of the J/ψ, this gives us $\sigma_t(J/\psi) \simeq 4$ mb. To fit the existing J/ψ suppression with the form (19), we assume that absorptive processes provide half of the total cross section, and that the initial density $n(\epsilon_o)$ of the constituents (including resonances) is 2.5 times smaller than that obtained from the observed pions [17–19]. The resulting J/ψ suppression is included in Fig. 10. Within the experimental errors, it agrees with the data as well as the deconfinement form (18). Clearly better statistics are needed.

Finally we want to point out that a comparison of data for J/ψ and ψ' production will provide a very useful tool to probe the nature of the primordial state and to distinguish between deconfinement and absorption.

We recall that in p-A collisions, the ratio of $(S_{\psi'}/S_{J/\psi})$ is in good approximation unity [23]. This means that if we observe deviations from unity, we are "automatically" looking at phenomena beyond whatever is responsible for the suppression in p-A collisions, whether this is a structure function modification, as claimed here, or absorption of the nascent J/ψ [27,36], colour independent comover interactions [11], parton energy degradation [37,38], or a combination of these. A deviation of $(S_{\psi'}/S_{J/\psi})$ from unity implies that the medium can distinguish between the two states, be it through the different dissociation energy densities ϵ_d, through the different absorption cross sections, or through the difference between the charmonium mass and open charm (the $D\bar{D}$ threshold). For the presently accessible x_F range in p-A collisions, the medium can appearently not do this, and whatever the origin of a possible $(S_{\psi'}/S_{J/\psi}) \neq 1$ might be, it would not be obtainable by extending p-A arguments. Let us then compare the effects of deconfinement and absorption on the ratio in question.

If colour deconfinement is the dominant suppression mechanism, then the ratio of ψ' to J/ψ production is

$$\frac{S_{\psi'}}{S_{J/\psi}}(\epsilon_o) = \frac{\Theta(1.9\epsilon_c - \epsilon_o)}{0.3 + 0.7(\epsilon_o/\epsilon_c)^{9/4}} + 0.3\ \Theta(\epsilon_o - 1.9\epsilon_c) \tag{19}$$

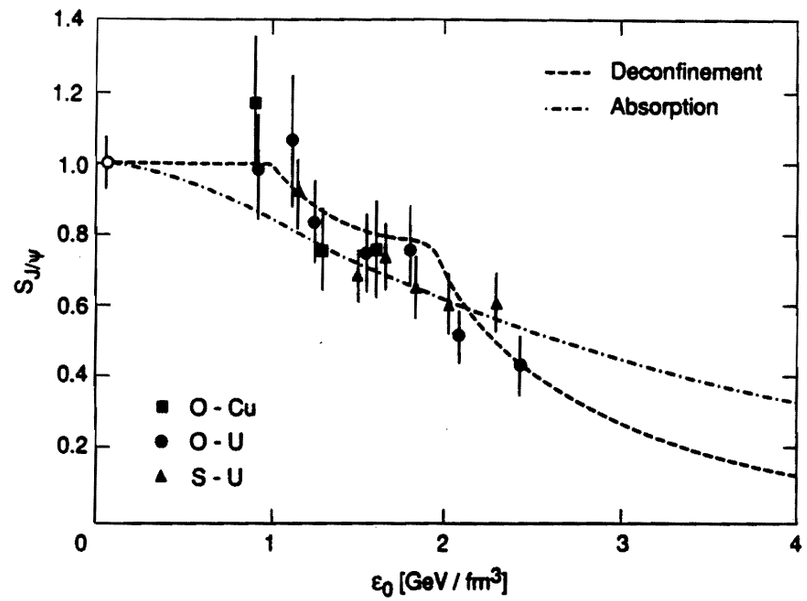

Fig. 10 Final state J/ψ suppression in nucleus-nucleus collisions as a function of the energy density [21], compared to the forms predicted by deconfinement (solid line) and absorption (dashed line). The data are normalised to the $p-U$ value, with errors as indicated on the open circle.

where we have used $\epsilon_{\psi'} \simeq \epsilon_c$ (see Table 1). The resulting behaviour is shown in Fig. 11: for $\epsilon_o \leq \epsilon_c$, we have $(S_{\psi'}/S_{J/\psi}) = 1$, and for all $\epsilon_o \geq \epsilon_\psi$, we get

$$\frac{S_{\psi'}}{S_{J/\psi}} = \frac{1}{0.3 + 0.7(\epsilon_\psi/\epsilon_c)^{9/4}} \simeq 0.3. \tag{20}$$

The transition between these two regimes is "softened" by the two-state origin of the J/ψ (and by quantum effects [30]); break-up collisions between the $c\bar{c}$ and partons in the plasma would also tend to soften the transition [39].

In the absorption scenario, on the other hand, both the J/ψ signal itself and the ratio of ψ' to J/ψ production should decrease exponentially with increasing energy density,

$$\frac{S_{\psi'}}{S_{J/\psi}}(\epsilon_o) \simeq \exp\{-\Delta\sigma\, n(\epsilon_o)\tau_o \ln[n(\epsilon_o)/n_f]\}\,, \tag{21}$$

where $\Delta\sigma = \sigma(\psi') - \sigma(J/\psi)$ denotes the difference between the absorption cross sections of the $\psi'$and the J/ψ. The corresponding behaviour of the $\psi'/(J/\psi)$ ratio, with the same input parameters as above and $\sigma_t(\psi') \simeq 11$ mb, is included in Fig. 11.

The essential qualitative difference thus is that the ratio of ψ' to J/ψ production in the colour deconfinement scenario becomes constant, once ϵ_o is above the relevant threshold densities, while in the absorption picture the ratio continues falling exponentially for all ϵ_o. In the energy density range from 2 to 6 GeV/fm^3, which should become accessible at the latest with RHIC, deconfinement gives a constant $S_{\psi'}/S_{J/\psi}$, while for absorption, this ratio falls by a factor four.

We thus see that absorption and deconfinement lead to quite different behaviour at higher energy densities, and this holds for any ratio of charmonium and bottonium states. A full $c\bar{c}/b\bar{b}$ spectral analysis in nuclear collisions therefore promises to be a viable tool to probe the primordial state of the dense matter produced in high energy nuclear collisions.

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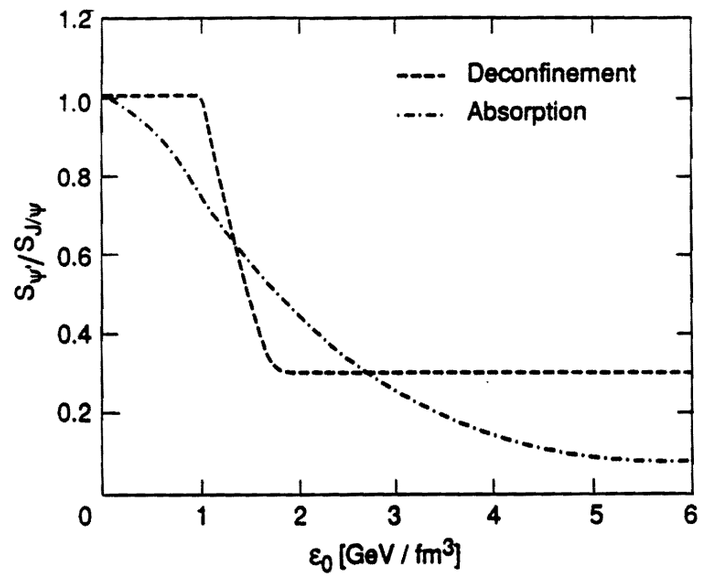

Fig. 11 The ratio of the survival functions $S_{\psi'}/S_{J/\psi}$ as a function of the energy density for deconfinement (dashed line) and absorption (dash-dotted line).

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