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Charmonium suppression by thermal dissociation and percolation

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Abstract. I discuss the charmonium suppression in deconfined medium by thermal dissociation and parton percolation. I point out the differences and show predictions for J/ψ suppression at different energy and/or for different interacting nuclei.

1 Introduction

The behavior of J/ψ mesons in a hot strongly interacting medium was proposed as a test for its confinement status [1]: it was argued that the J/ψ, due to its small size and strong binding energy, cannot break up as a consequence of interactions with normal hadrons, while in a deconfined medium, the color screening dissolves the $c\bar{c}$ bond.

After this proposal, the study of the charmonium suppression in heavy ion collisions has aroused a lot of interest. With a careful and extensive analysis of the experimental data for different interacting systems (from proton–proton to proton–nucleus and nucleus–nucleus collisions) it has been possible to observe a "normal" suppression of the J/ψ meson, presumably due to the absorption of the preresonant $c\bar{c}$ state in the nuclear medium, in all interactions up to S–U collisions and peripheral Pb–Pb. In central Pb–Pb collisions a stronger suppression is observed [2], suggesting that a new suppression mechanism is at work in these events.

The aim of this work is to present an interpretation of the "anomalous" J/ψ suppression as a consequence of a deconfinement transition, in the framework of two models, thermal dissociation and parton percolation. I will focus on $c\bar{c}$ states, but all the considerations presented here apply to the $b\bar{b}$ case as well.

A few general remarks are necessary before presenting the theoretical models.

It is known from proton–nucleon and pion–nucleon [4] (and, more recently, from electron–proton [5]) interactions that a large fraction (about 30%–40%) of the observed J/ψ's are the decay products of higher excited states of $c\bar{c}$ pairs (ψ', χ). Since the life-time of these quarkonium states is much larger than the typical life-time of the medium which is produced in the nucleus–nucleus collisions, they decay presumably in the vacuum. Therefore this medium (either hadronic gas or quark–gluon plasma) sees not only the ground state quarkonium, but also the different excited states, which have different properties (size, binding energies) and different behavior: a smaller binding energy (and, consequently, a larger radius) requires a lower dissolution temperature, in the case of a deconfined medium; on the other hand, for the case of a hadronic system, a weakly bound quarkonium state has a larger break-up cross-section for interactions with the other particles.

Therefore the final J/ψ survival probability, to be compared to the experimental data, has to be calculated as an average over the different components, each of them weighted with the corresponding fraction of contribution to the observed J/ψ's in the final state:

$$S_{J/\psi} = f_{J/\psi} S^{\rm dir}_{J/\psi} + f_\chi S^{\rm dir}_\chi + f_{\psi'} S^{\rm dir}_{\psi'}. \tag{1}$$

This fact has a very important consequence on the pattern of the J/ψ suppression as a function of the centrality and of the energy collisions and it must be considered for a careful comparison to experimental data.

2 Thermal dissociation

In sufficiently hot deconfined matter, color screening dissolves the binding of the quark–antiquark pair, and a stronger binding energy requires a higher temperature to be dissolved. On a microscopic level, it was argued that only a hot medium provides sufficiently hard gluons to dissociate the quark–antiquark bound state, and this again implies a dissociation hierarchy as a function of the binding energy. Another possible mechanism is the decay into open charm mesons due to in-medium modification of mesonic masses [3].

Implicit, in this approach, is the assumption that the medium probed by the quarkonium state is in thermal equilibrium. Lattice studies show that the transition between confined and deconfinement medium occurs at the critical temperature $T_{\rm c} \simeq 150$–$200\,$MeV.

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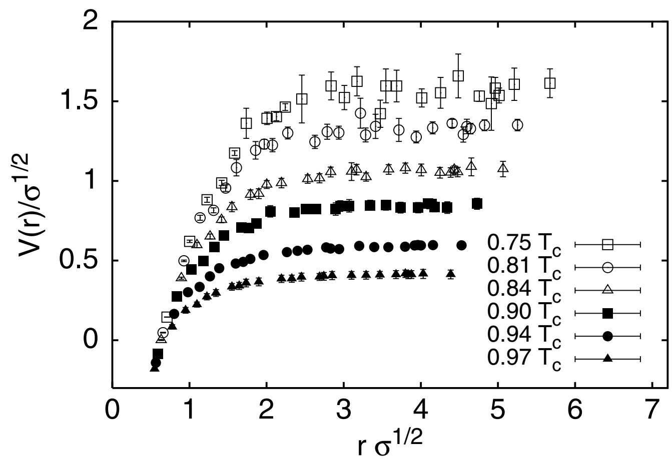

Fig. 1. The heavy quark potential below T_c at different temperatures [8]

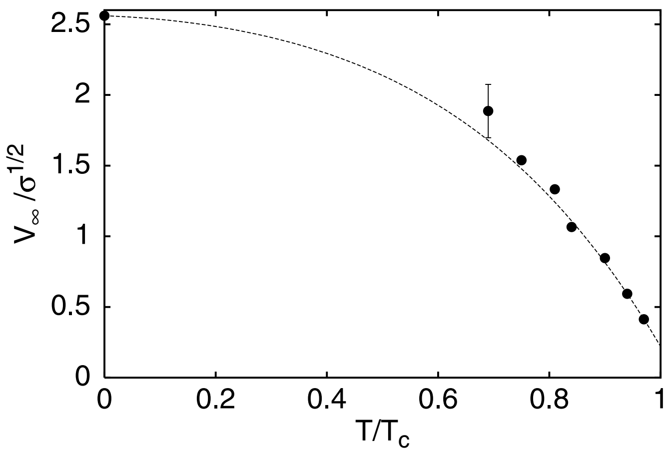

Fig. 2. The asymptotic value of the heavy quark potential below deconfinement. Points are the lattice results [8], the line is a fit [3]

2.1 Quarkonium dissociation below deconfinement

The large values of the charm (and bottom) quark mass allows potential theory to describe quite accurately the quarkonium spectroscopy [6]. At zero temperature, a phenomenologically successful description is given by the Cornell potential [7]:

$$V(r) = \frac{a}{r} + \sigma r \ . \qquad (2)$$

More generally, a precise form for the potential is obtained, at any temperature, from first principle QCD with lattice calculations [8]. The potential is related to the Polyakov loop $L(r)$, calculable on lattice, by

$$V(r,T) = -\log\langle L(r)L^\dagger(0)\rangle + C \ , \qquad (3)$$

where C is a normalization constant. In absence of dynamical quarks, $V(r,T)$ is linearly rising for large r; if dynamical quarks are present, one expects that at a distance r such that $V(r,T)$ is equal to twice the mass of a light quark, the string between the heavy quarks breaks, a pair of light quarks $q\bar{q}$ is produced, and the quarkonium decays in a pair of charmed mesons ($Q\bar{Q} \to Q\bar{q} + \bar{Q}q$).

At very short distances medium effects are not important; therefore, for very small r, the potential should always be of the form of (2). In [3] the constant C was determined by imposing that, at any temperature below T_c, the potential given by lattice calculation reproduce the Cornell form for $r \to 0$. The results are shown in Fig. 1. It is evident that the asymptotic value $V_\infty(T)$, for $r \to \infty$, decreases rapidly approaching the critical temperature, as shown more clearly in Fig. 2.

At any temperature T, a functional form of the heavy quark potential can be obtained by fitting the lattice results: this potential can then be used in the Schrödinger equation to calculate the binding energy (i.e. the mass) of the ground state and of the first excited states of the $Q\bar{Q}$ pair. The results for the J/ψ (1s), χ (1p) and ψ' (2s) states are shown in Fig. 3 (from [3]) as a function of the temperature. In the same figure the solid line represents

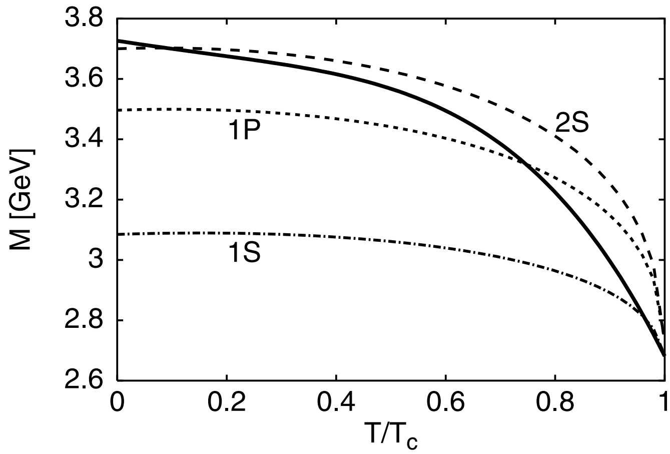

Fig. 3. The open charm threshold (solid line) and the charmonia masses (dashed lines) as a function of the temperature, below T_c [3]

$V_\infty(T)$, as obtained from the fit shown in Fig. 1: when the mass of a given $c\bar{c}$ bound state is larger than $V_\infty(T)$, the dissociation into a pair of charmed mesons is energetically favored. One therefore concludes that the ψ' is so weakly bound that it dissolves already at low temperatures, while the χ requires $T \simeq 0.75T_c$; the J/ψ, on the contrary, seems to survive up to the critical temperature.

2.2 Quarkonium dissociation in a deconfined medium

Above the critical temperature T_c one defines the constant C in such a way that the color average potential vanishes at large distances, as shown in Fig. 4. It is also evident that the value $r_0(T)$ at which $V(r,T)$ vanishes decreases rapidly with increasing T (thick line of Fig. 5).

The color average potential calculated on the lattice is the thermal average of the color singlet and color octet contributions:

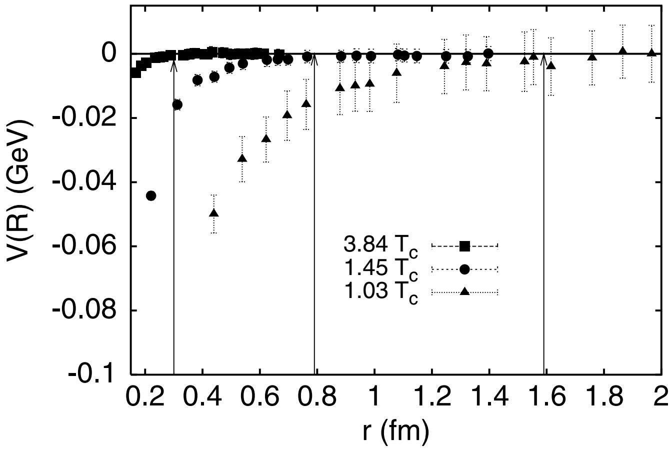

Fig. 4. The color average potential calculated on the lattice above T_c [8]

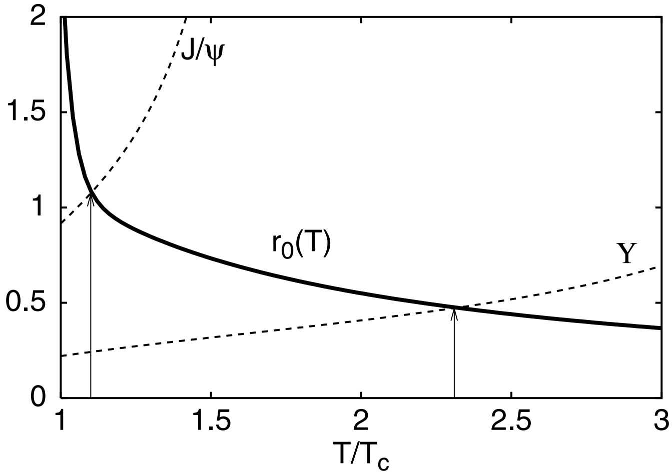

Fig. 5. The effective radii of the J/ψ and Υ states [3]

$$\exp\left\{-\frac{V(r,T)}{T}\right\} = \frac{1}{9}\exp\left\{-\frac{V_1(r,T)}{T}\right\} + \frac{8}{9}\exp\left\{-\frac{V_8(r,T)}{T}\right\}. \quad (4)$$

Assuming for the singlet and the octet potential, respectively, the following forms (valid at leading order in perturbation theory):

$$V_1(r,T) = -\frac{4}{3}\,\frac{\alpha(T)}{r}\,\exp\{-\mu(T)r\}, \quad (5)$$

$$V_8(r,T) = \frac{c(T)}{6}\,\frac{\alpha(T)}{r}\,\exp\{-\mu(T)r\}, \quad (6)$$

one can fit the lattice results and extract the singlet potential. The temperature-dependent coefficient $c(T)$ has been included to take into account the strong non-perturbative effects just above T_c; for $T \ll T_c$ one expects that $c(T) \simeq 1$ and the usual $1:8$ ratio between singlet and octet components is recovered.

To see whether quarkonium bound states can exist above the critical temperature, one solves the Schrödinger

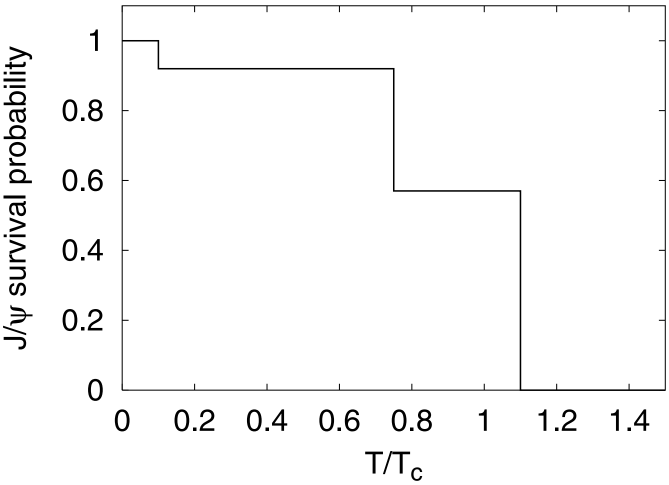

Fig. 6. The J/ψ suppression pattern as a function of the temperature, assuming 57% of directly produced J/ψ's, 35% and 8% of feed-down from χ_c's and ψ''s

equation with the V_1 potential (V_8 is repulsive):

$$\left\{2m + \frac{1}{m}\Delta + V_1(r)\right\}\psi_i = M_i\psi_i\ . \quad (7)$$

A relevant piece of information that is obtained from (7) is the radius r_i of the quarkonium bound state: when $r_i(T)$ is larger than the screening distance of the medium $1/\mu(T)$ the corresponding bound state dissolves. Figure 5 shows that for the J/ψ this happens at about $1.1T_c$, while the Υ survives up to temperatures well above $2T_c$.

The consequent J/ψ suppression pattern, assuming that 35% and 8% of the observed J/ψ's come from χ and ψ' decays respectively, is similar to what shown in Fig. 6 as a function of the temperature. One should assume some model to relate the temperature to some experimental observable, in order to study the suppression as a function of the centrality of the collision, and what is obtained is a "two-step" pattern, which will be smoothed by fluctuation of the number of participants (or any other relevant observable) with the impact parameter. If the colliding nuclei are not heavy enough, it is possible that only the first threshold is reached; the second step therefore will not be present. But, on the other hand, if the incident energy of the collision is very high, it is also possible that the two steps occur for very peripheral collisions and very close to each other (i.e. in a small range of impact parameter) so that they cannot be easily resolved.

3 Parton percolation

Hadrons are made by partons. When two or more hadrons overlap, their partons interact. In a normal hadron–hadron interaction, the overlapping phase is so short in time that the partons separate again into hadrons before reaching an equilibrium condition. On the other hand, in a nucleus–nucleus collision, the number of interacting hadrons is so high that their partons can interact several times, they therefore lose their "identity", so that they do not belong anymore to a particular hadron but form a big cluster of deconfined medium: the quark–gluon plasma (QGP).

The percolation theory is a mathematical tool which studies how simple objects form clusters; there are applications of the percolation idea in many physical problems and the deconfinement transition in strongly interacting matter is one of them. For instance, in [9] hadrons interact by exchanging color strings. When many hadrons interact simultaneously in a small space-time region, these strings overlap and, when their density reaches a critical value, they percolate. The model of hadron interaction based on color string exchange (below the percolation threshold) is able to reproduce many features of experimental data (see references in [9]). This model therefore interpolates nicely from interactions in a normal hadronic medium and deconfined matter.

A similar approach was followed in [10]: a model of string fusion and percolation is used to describe several experimental observables, including J/ψ suppression.

The work of [11–13] is essentially focused on J/ψ suppression by parton percolation, inspired by lattice results where it was shown that the deconfined transition in SU(2) gauge theory can be described by percolation of Polyakov loops [14].

Parton percolation is an essential prerequisite for QGP formation. It should be noted that thermal equilibrium is not required and this is a main difference with respect to the approach described in Sect. 2.

3.1 Basic concepts of percolation

Consider a two-dimensional circular region of radius R on which N small discs of radius $r \ll R$ are randomly distributed. When two or more small discs overlap, they form clusters. The cluster size increases with increasing $n = N/\pi r^2$, the average disc density, and this cluster formation shows critical behavior: in the limit $N \to \infty$ and $R \to \infty$, with n finite, the cluster size diverges at the critical density n_c, given by

$$n_c = \frac{\nu_c}{\pi r^2}, \tag{8}$$

where the critical coefficient $\nu_c \simeq 1.13$ has been determined by numerical calculations. In a finite system, the cluster size increases very rapidly as a function of the filling factor $\nu = n\pi r^2 = Nr^2/R^2$, at the percolation onset, as shown in Fig. 7, where the location of the critical filling factor ν_c is indicated by an arrow.

The percolation onset is defined as the point at which the growth of the cluster is more rapid [11], i.e. where its derivative with respect to ν picks (dashed line of Fig. 7). This definition is appropriate for a finite system and it is an obvious generalization of the more traditional definition for an infinite region.

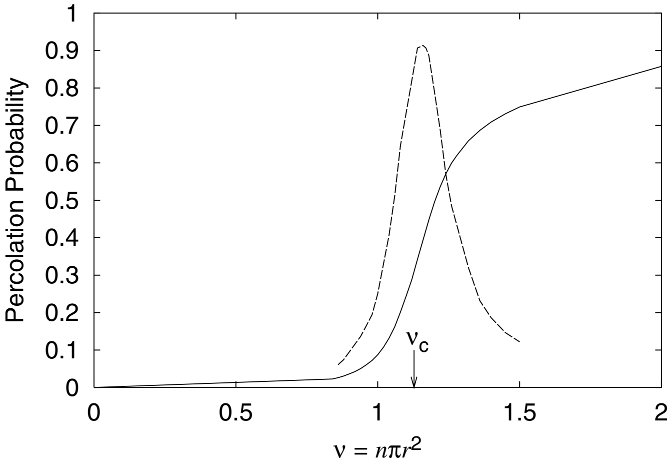

Fig. 7. Cluster size (solid line) and its derivative as a function of the filling factor, from [12]

3.2 Percolation in high energy nuclear collisions

The above considerations apply to a nucleus–nucleus collision if one makes the following assumptions: the overlapping objects, forming clusters, are colored partons and the two-dimensional space where they are distributed is the transverse plane projection of the overlapping region of the two incident nuclei. If the collision is central (impact parameter $b = 0$), the overlapping region is circular, as in the example discussed above. In the case of a collision at $b > 0$, the overlapping region will be almond-shaped and the geometry of the collision will be a little more complicated, but all the considerations presented here are valid. A parton cluster represents a region in which color charges can move freely: if it extends over the entire region, one has, by definition, color deconfinement.

In high energy nuclear collisions there is one additional difficulty: the partons are emitted by the interacting nucleons, therefore their distribution is not uniform in the transverse plane, but rather reflects the original distribution of the initial nucleons, with a higher concentration in the center than near the surface of the nucleus. It is possible that in a given collision, only the most central region of the produced medium is dense and hot enough to allow for the deconfinement transition. In this situation a local definition of percolation, as the one used in [12], is more appropriate. At the onset of percolation, the parton density m in the largest cluster is slightly larger than the overall average density n; in fact, numerical studies show that at the critical density n_c, the density in the largest cluster reaches $m \equiv m_c = \eta_c/\pi r^2 \simeq 1.72/\pi r^2$: this value provides the local percolation condition: if the parton density at a certain point exceeds this value, the corresponding cluster percolates and the medium, in the region occupied by it, is deconfined. The relevance of this local percolation condition for the J/ψ suppression in nuclear collisions is evident: the J/ψ meson is very small and therefore it is sensitive to the properties of a small spatial region and, moreover, it allows for the deconfined matter to be produced only in a limited part of the produced medium.

Having specified the objects which form clusters, the partons, one has to describe their distribution in space to apply the percolation idea to nuclear collisions. Since these partons are emitted by the incident nucleons as a consequence of the strong interactions during the first stages of the nuclear collision, it is reasonable to assume that their number and spatial distribution is determined by the participating nucleons from which they originate, as done in [12, 13]. Therefore the density of partons is given by the product of the participating nucleon density $n_s(b, A)$ (which depends on the nucleus A and on the impact parameter of the collision) and the number of partons per nucleon $\mathrm{d}N_q(x, Q^2)/\mathrm{d}y$ (the parton distribution function, known from deep inelastic scattering experiments). The fraction x of the nucleon momentum carried by the parton is related to the incident energy $\sqrt{s}$ by $x = (k_\mathrm{T}/\sqrt{s})$ (at midrapidity); k_T is the average transverse momentum of the parton and it is inversely proportional to its transverse size. The global percolation condition is then given by

$$n_s(A, b) \left(\frac{\mathrm{d}N_q(x, Q_\mathrm{c}^2)}{\mathrm{d}y} \right) \Bigg|_{x=Q_\mathrm{c}/\sqrt{s}} = \frac{\nu_\mathrm{c}}{(\pi/Q_\mathrm{c}^2)} \ . \qquad (9)$$

Equation (9) can be solved numerically to obtain for what value of A at a given $\sqrt{s}$ percolation sets in (or, vice versa, for which energy percolation can occur in collisions of nuclei with a given atomic number A) and the value of Q_c at the percolation point.

From a local point of view, the density of the largest cluster at the percolation point gives the local percolation condition: $m_\mathrm{c}(A, b) = \eta_\mathrm{c}/(\pi/Q_\mathrm{c}^2)$. With these ingredients, the authors of [13] find that the critical cluster density is reached, in Pb–Pb collisions at SPS, at $b \simeq 8$ fm and the corresponding value of the average transverse momentum of the partons is $Q_\mathrm{c} \simeq 0.7$ GeV. The scales of the charmonium states χ_c and ψ', given by the inverse of their radii calculated in potential theory, are about 0.6 and 0.5 GeV respectively: they are therefore dissociated at the onset of percolation. On the other hand, directly produced J/ψ's have smaller radii; therefore the average transverse momentum of the deconfined partons must be at least 0.9–1.0 GeV to resolve them, so only a denser medium, produced in more central collisions, can dissociate them.

The J/ψ survival probability in nuclear collisions is then obtained, from the above considerations, by assuming that about 40% of the J/ψ's (those coming from χ_c and ψ' decays) produced inside the percolating cluster disappear at the onset of percolation (those formed outside the cluster, i.e. near the surface, are not affected). The remaining 60% of J/ψ's (the directly produced ones) survive until a cluster of hard enough partons is produced ($b \simeq 3$–4 fm). For a realistic comparison to the experimental data one has, of course, to take into account impact parameter fluctuations (see, for instance, [15]). The result is shown in Fig. 8 [13], compared to the experimental data provided by the NA50 Collaboration; one sees that the percolation model reproduces quite well the features of the experimental data, in particular the "two-step" pattern.

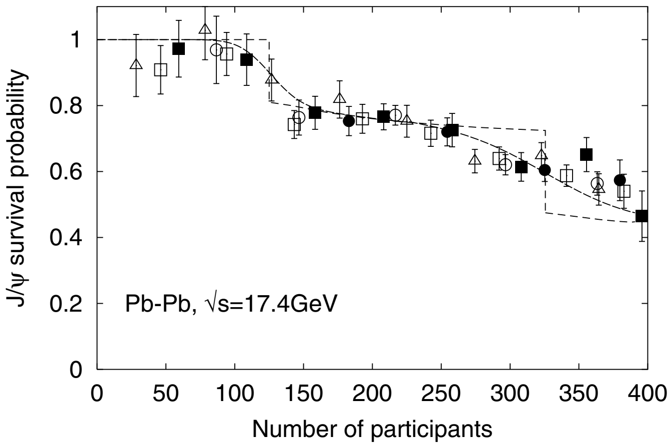

Fig. 8. J/ψ survival probability as a function of the number of participants, in Pb–Pb collisions at SPS. Experimental data from NA50 Collaboration

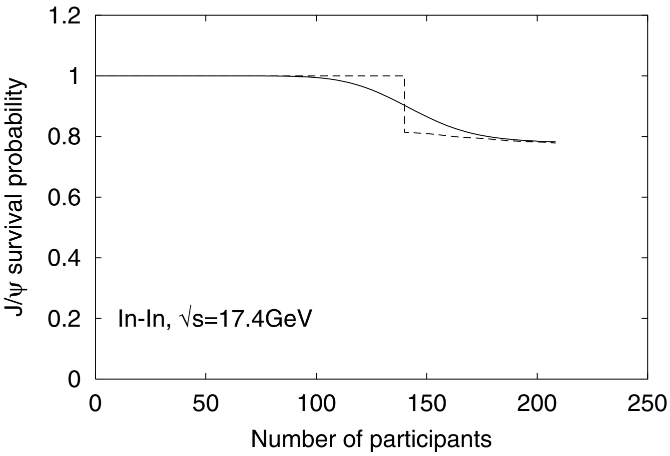

Fig. 9. J/ψ survival probability in In–In collisions at SPS

The same model can be used to predict the survival probability as function of the centrality in In–In collisions at SPS. One finds, in this case, that the percolation onset is reached in semi-central collisions ($N_\mathrm{part} \simeq 140$), but the threshold for the dissociation of directly produced J/ψ's is never reached. One therefore expects, in this model, to observe a suppression pattern similar to that presented in Fig. 9, with one step only.

At higher energies, things can change. For instance at RHIC energy the parton density, as given by the parton distribution functions, is so high that at the percolation onset the average transverse parton momentum is hard enough to resolve all the charmonium states, including the directly produced J/ψ's. The suppression pattern has again an unique step, but deeper than the previous case (where only higher excited charmonium states were affected) as shown in Fig. 10.

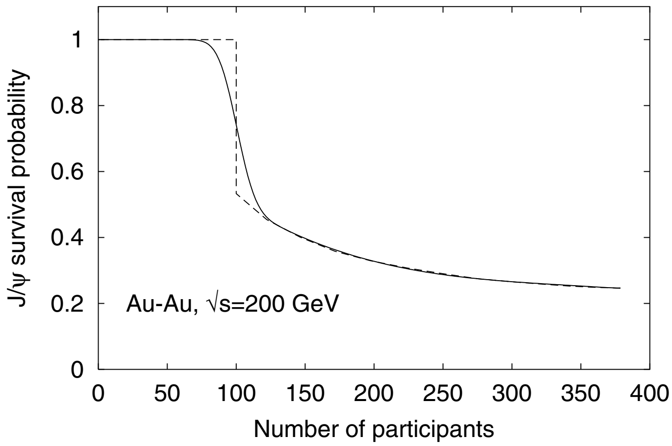

Fig. 10. J/ψ survival probability in Au–Au collisions at RHIC

4 Conclusions

The two theoretical models presented here give a good description of the experimental data on J/ψ suppression in Pb–Pb collisions at the SPS. The main difference is that the dissociation model requires a thermalized medium, whereas the percolation model does not. The thermal dissociation model seems, presently, disfavored by lattice results where the J/ψ is found to survive up to very high temperatures (2–$2.5T_{\mathrm{c}}$), very difficult to accommodate in this approach; but the lattice results [16], on the other hand, are not yet conclusive since the widths of the charmonium states are not calculated so it is difficult to deduce, from them, the real behavior of a J/ψ in a deconfined medium.

Another difference in the two models is evident in the predictions for higher energies: the thermal dissociation model always predicts a sequential suppression of the different quarkonium states, giving rise to a "two-step" pattern (which can be more or less evident, depending on the range of the impact parameter in which the two thresholds occur), whereas the percolation approach can give, as in the case of Au–Au at RHIC, a qualitatively different suppression as a function of the centrality.

It should be noted that the "two-step" pattern observed in the experimental data at SPS can be explained by other mechanisms. In particular, it has been argued that the second "drop", in central Pb–Pb collisions, is due to strong energy density fluctuations [17]: if this is the case, the same pattern, i.e. the two steps, should be observed in In–In collisions as well, where the percolation model, on the contrary, predicts only one step.

I did not discuss here the charmonium suppression by the percolation models of [9,10]: the agreement between theory and experiment may seem poorer than the results presented above [13], but one should keep in mind that in these works the feed-down from χ_c and ψ' is neglected and, more importantly, the free parameters are tuned to reproduce other observables of the nuclear interactions: in view of that, the agreement obtained in the J/ψ case is remarkable and it suggests that the percolation approach can be extended to give a more general description of nuclear collisions.

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