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The evolution of the quark-gluon plasma

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Abstract. We survey the different stages in the evolution of the quark-gluon plasma expected to be produced in high energy nuclear collisions, and we discuss some specific experimental probes for each.

1. Introduction

When we study strongly interacting matter in the framework of statistical quantum chromodynamics, we deal with ensembles of time-independent systems. Adiabatically, quark matter becomes hadronic matter, or vice versa. Our systems are in equilibrium to begin with, and they are transformed, very slowly and by slight of changes in external parameters, into other equilibrium systems. Nevertheless, we hope to test the essential predictions of these studies, quark deconfinement and chiral symmetry restoration, through nuclear collisions, in which our initial system, consisting of two colliding nuclei, is certainly not in any way thermal, and in which whatever matter is eventually produced very rapidly expands and freezes out into non-interacting hadrons. While the basis of what we try to study indeed is given by statistical QCD, our tools must clearly be adopted to the rapidly evolving world accessible to us in high energy nuclear collisions. Whatever we measure will reflect the properties of the system in some stage of its evolution, and we have to learn to associate specific observables and specific stages of the system we are studying. In this survey, I will try to give some examples of how to probe the evolution of the quark-gluon plasma, from two colliding parton beams to the observation of identified hadrons.

Over the past years, three stages in high energy nuclear collisions have attracted particular attention. In the period immediately after the collision, multiple scattering between the primary and secondary constituents is supposed to lead to a rapid generation of entropy and thus eventually to thermalisation. What happens at this stage is of crucial relevance for the subsequent properties of the system, and much important work on this topic has appeared during the past year, based in particular on event generators using parton dynamics [1]. The only tools we have to check this stage directly are hard early phenomena suffering little or no effect by the further evolution of the system after the pre-equilibrium stage, such as high mass Drell-Yan dileptons or $D\bar{D}$ pairs. At the other end of the evolution time scale are hadrons presumably emitted from a thermal source. They will reflect the state of the system at the time they were formed – i.e., they will tell us about the hadronisation of the quark-gluon plasma [2]. If nuclear collisions do indeed lead to a thermalised system of quarks and gluons, this hadronisation stage, the final state in the evolution, will be independent of the initial thermal state, the primordial plasma. In other words, the *local* hadronisation properties of a plasma initially at a temperature of one GeV should be no different from those of a plasma of two hundred MeV [3]. How can we then probe the primordial plasma, the system we expect to be hotter at RHIC than at the SPS, and much hotter still at the LHC? The tools proposed for this purpose include quarkonium spectra [4], high mass thermal dileptons and hard thermal photons [5], thermal charm [6], and hard jets [7]. Of these, up to now only J/ψ suppression – predicted as a consequence of a hot primordial quark-gluon plasma – was actually observed; the analysis of charmonium spectra is therefore of particular interest.

In the following, I will address some specific aspects of probing each of the three evolution stages. Section 2 will be devoted to pre-equilibrium and thermalisation. In Section 3, quarkonium production will be considered as a probe of the early hot and dense thermal stages in nuclear collisions. Section 4 will then deal with what spectra of thermal hadrons can tell us about the quark-hadron transition and the freeze-out into free hadrons.

2. Pre-equilibrium and thermalisation

Thermal equilibrium can be attained by multiple scattering among the constituents of the incident nuclei and the secondaries produced in these interactions. The primary interactions must trigger a cascade of more and more interacting secondaries, tertiaries, and so on, leading to a rapid entropy increase and with it a (local) loss of memory of the initial state. Computer simulations of such cascades provide estimates of the time t_0 needed to reach local equilibrium, and they tell us what the thermal conditions were at that time. They specify the initial energy density ϵ_0, the initial entropy density s_0, and they can tell us to what extent there is "chemical" equilibrium among the different constituent types, whether hadrons or partons.

In such simulations, a considerable number of input aspects has to be specified in order to arrive at definite predictions.

- What is the nature of the constituents? They can be hadrons, constituent quarks, or the partons (massless quarks and gluons) of perturbative QCD.

- What is the collision dynamics? We can start from a phenomenological picture of hadronic interactions (ARC [8], dual parton model [9]), from string interactions (Venus [10], RQMD [11] Spacer [12]), or from perturbative QCD for partonic interactions, with soft additions both for low P_T processes and for the hadronisation of the partons (Hijing [13], parton cascade model [14]).

- What is the evolution pattern for the collision? Here we can consider independent successive binary collisions, molecular dynamics schemes or a more detailed transport theory description.

- What collective interaction effects should be included? Since all evolution patterns considered are approximative, we have to make sure that such collective phenomena as Pauli blocking, the Landau-Pomeranchuk effect, parton shadowing and jet quenching are correctly taken into account.

It is thus clear that the results of such microscopic evolution studies are necessarily quite model dependent. Therefore an obvious requirement is that before making predictions for nuclear collisions, one has to tune the particular generator to reproduce the available data from hadron-hadron and hadron-nucleus collisions. Beyond this one may hope that in the high energy limit things become less ambiguous: eventually the basic input should be an internetted parton cascade, based on an underlying description in terms of perturbative QCD. The past year has brought a number of interesting studies in this direction. Here we can only summarize their results; for details, we refer to [15] [16].

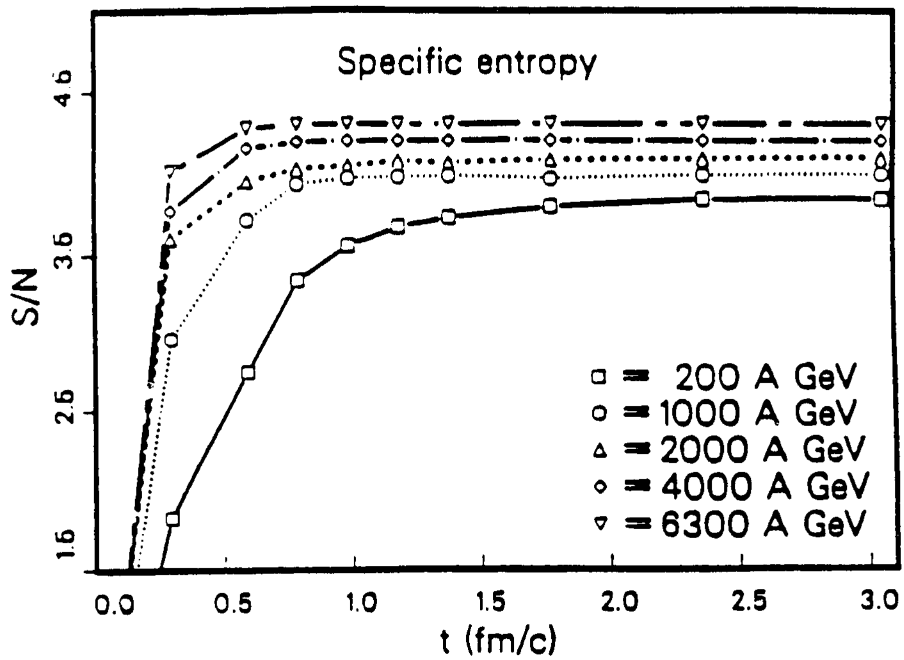

Figure 1. Entropy per parton as function of time after the collision, for different incident CMS energies $\sqrt{s}/A$; from [16].

The specific entropy (entropy per secondary) increases very rapidly at high energies [16], as seen in Fig. 1. While at RHIC energy it is still compatible with the one fermi generally assumed to be the equilibration time at present energies, we approach the expected value $S/N = 4$ in less than 0.5 fm for LHC energy. This rapid equilibration is also reflected in increasingly shorter relaxation times (Fig. 2). As a consequence, the initial temperatures become higher (Fig. 3), since we can now start considering the system as thermal at earlier times.

On the other hand, these high energy considerations do not lead to a corresponding rapid onset of chemical equilibrium. As had been pointed out quite some

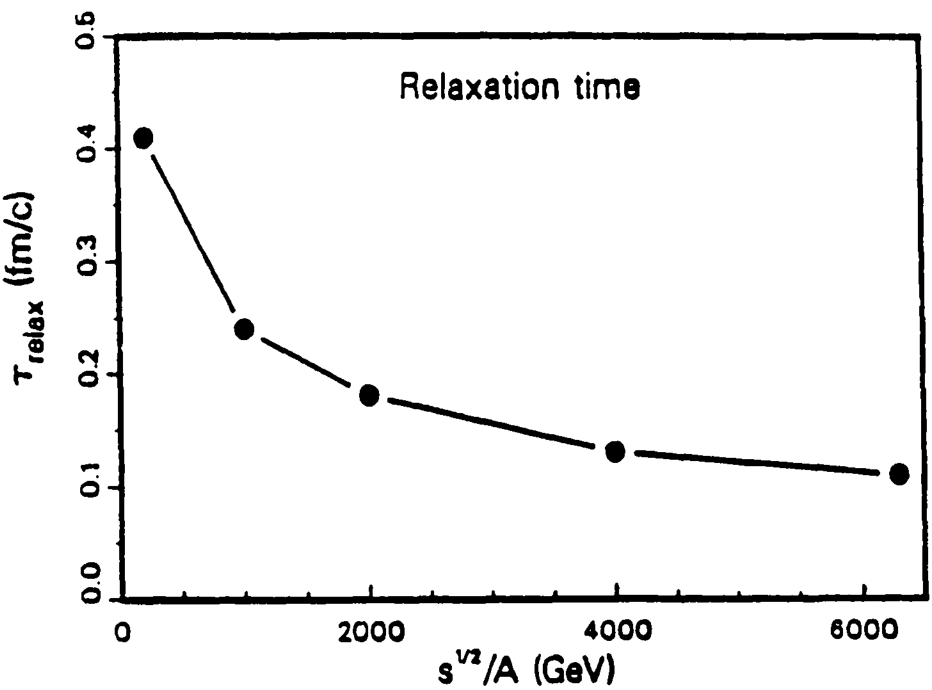

Figure 2. Relaxation time as function of the beam energy $\sqrt{s}/A$; from [16].

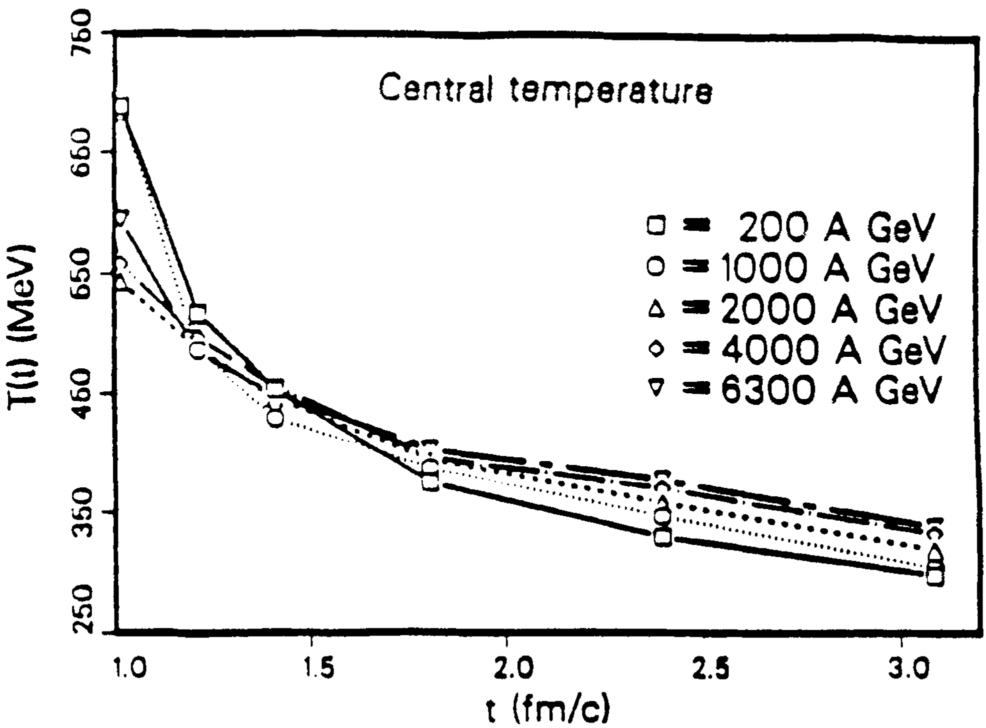

Figure 3. Central temperature as function of time after the collision, for different incident CMS energies $\sqrt{s}/A$; from [16].

time ago [17], the gluon interaction cross sections, including those leading to further gluon production, are considerably larger than their quark counterparts. This results in a perturbative QCD cascade with a very rapid gluon thermalisation, so that the entropy of the gluons soon reaches its equilibrium value. The entropy of the quarks, although it also increases quickly, appears to remain at about 60 - 70% of its equilibrium value for quite some time (Fig. 4). This behaviour has therefore been called the "hot gluon scenario" [6]. We should note, however, that it does not make the early thermal system a *gluon* plasma. A quark-gluon plasma (with 2.5 light quark flavours) in both thermal and chemical equilibrium would contain quarks and gluons in a ratio of about 1.6:1; in the hot gluon scenario we have instead a ratio of about 1:1, so that there are still as many quarks as there are gluons.

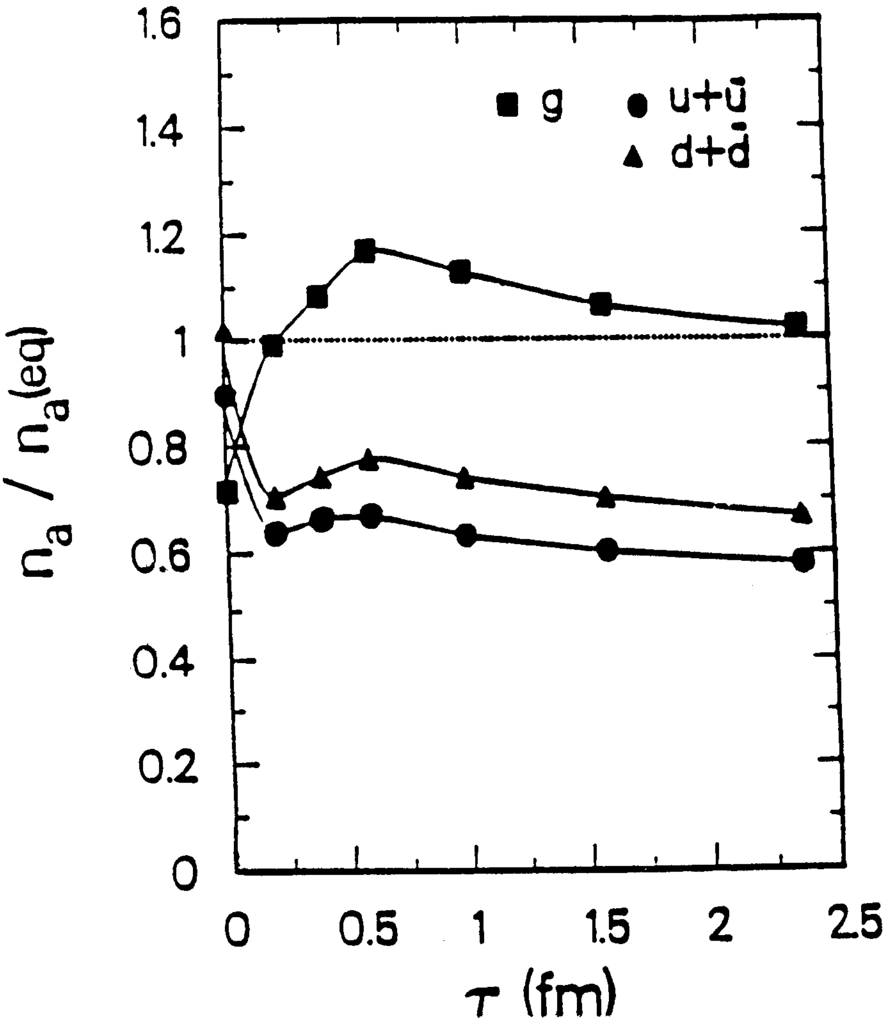

Figure 4. Parton densities, normalised to equilibrium values, for different parton species, as functions of the proper time after the collision; from [16].

The basis of these considerations is a dominance of perturbative QCD processes for the overall behaviour of the system at RHIC and in particular at LHC energies.

How can we test if this is correct? Two basic predictions of the partonic pre-equilibration stage are enhanced open charm [18] and enhanced high mass dilepton production [19]. The former arises because the abundant presence of hard secondary (and tertiary, etc.) gluons, the second since the fewer quarks are therefore harder. Since there are more hard gluons than quarks, open charm production will, however, be enhanced more than that of high mass dileptons.

Here two caveats seem appropriate. First of all, we have to say with respect to what the enhancement is to be measured. With decreasing fractional parton momentum $x \sim (M/\sqrt{s})$, the gluon distributions in *hadrons* are expected to grow more strongly than those for quarks, so that increasing energies will presumably lead to an increase of open charm relative to Drell-Yan productions even in $p-p$ collisions. Hence we will always need hadron-hadron collision data as reference to define the predicted pre-equilibrium enhancement in nucleus-nucleus collisions. Secondly, the high energy predictions for nuclear collisions have not yet been extrapolated down to lower energies where there is data to compare them to. This is obviously not trivial, since at lower energies the soft components are crucial. As long as their role is not checked specifically for the noted processes (open charm, Drell-Yan production), it is not clear how reliable the high energy predictions really are.

We summarize this section by noting that the production of open charm and beauty as well as that of high mass dileptons should provide good tools to probe the pre-equilibrium stage and the onset of thermalisation.

3. The primordial plasma

We now turn to the question of how we can verify if the systems produced in high energy nucleus-nucleus collisions were in their early stages indeed deconfined states of matter. As already noted, for this we need probes which are formed in the primordial stage, and only then, in order to avoid the subsequent "destruction of memory" in a thermal evolution. There are essentially three probes of this type which have so far been considered: quarkonium states (J/ψ, ψ', Υ, Υ'), real or virtual photons (dileptons), and hard jets. We shall here concentrate on the first of these three, since it constitutes the only predicted signature which has in fact been observed experimentally, since it has been studied most extensively (including, in particular, also competing alternative mechanisms), and finally since it illustrates the kind of analysis necessary before we can arrive at any conclusions. The topic of this section thus is quarkonium production as probe of the early quark-gluon plasma; to have a definite process, we shall talk about charmonium (J/ψ and ψ'), but everything is applicable to bottonium production as well.

The basic physics of the probe is quite simple. Charmonium production in hadronic collisions occurs in two stages. We have first a hard process, the fusion of two gluons to form a (coloured) $c\bar{c}$ pair, and then subsequently soft colour neutralisation and resonance binding. The first process, presumably described in terms of perturbative QCD, takes only about $M_c^{-1} \simeq 0.1$ fm or less; the second needs about 0.3 - 0.5 fm [20] [21]. If nuclear collisions lead to the formation of a deconfined medium, then the second step is inhibited in such collisions, the c and the $\bar{c}$ fly apart and eventually form open charm D mesons. In contrast, the overall Drell-

Yan rate is essentially unaffected by the presence of the plasma. In the dilepton spectrum from high energy nuclear collisions, we thus expect a suppression of the J/ψ and ψ' signals relative to the Drell-Yan continuum, if there was quark-gluon plasma production [4]. This suppression should manifest itself in a smaller signal-to-continuum ratio in nucleus-nucleus collisions when compared to hadron-hadron and hadron-nucleus collisions, and it should become stronger with increasing energy density in nucleus-nucleus collisions. Both these effects are in fact observed [22], and we thus have to interpret these observations.

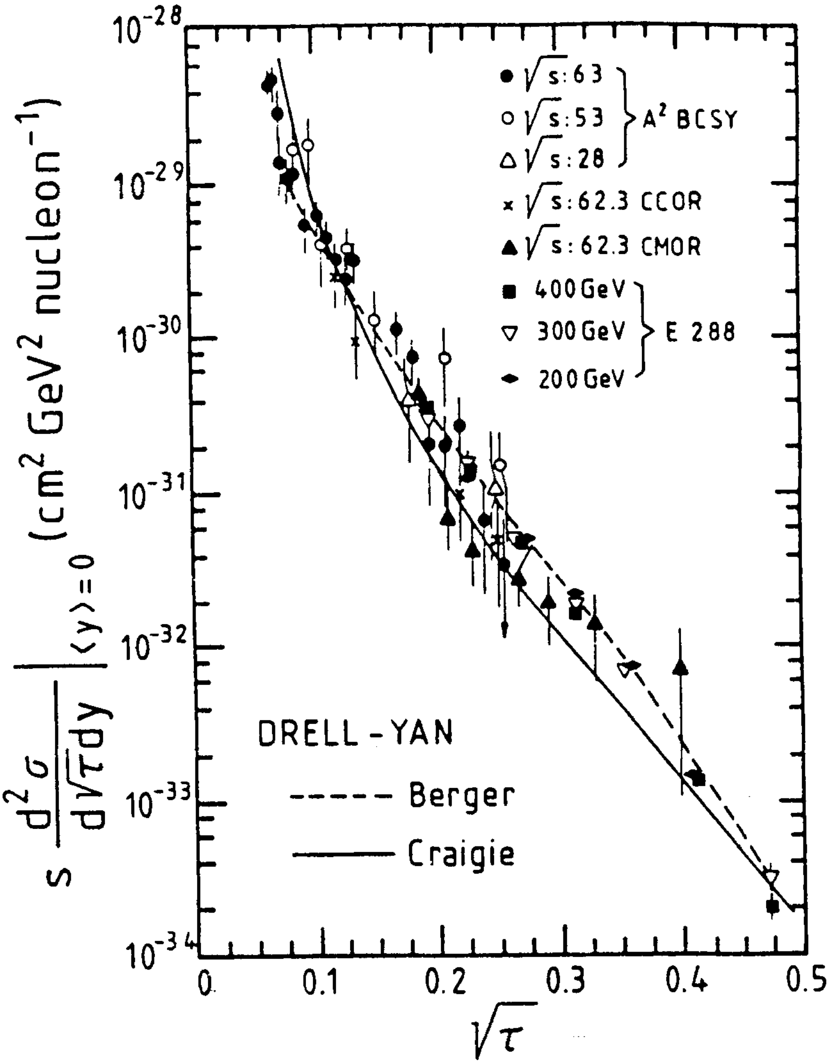

Figure 5. Drell-Yan rates in $p-p$ collisions [45], compared to calculations using quark structure functions (dashed line) [46] and to a fit [47].

Our analysis of experimental results will proceed in several steps. To start with, we show that the overall charmonium and Drell-Yan production rates in hadron-hadron collisions are basically understood, the first in terms of gluon fusion, the

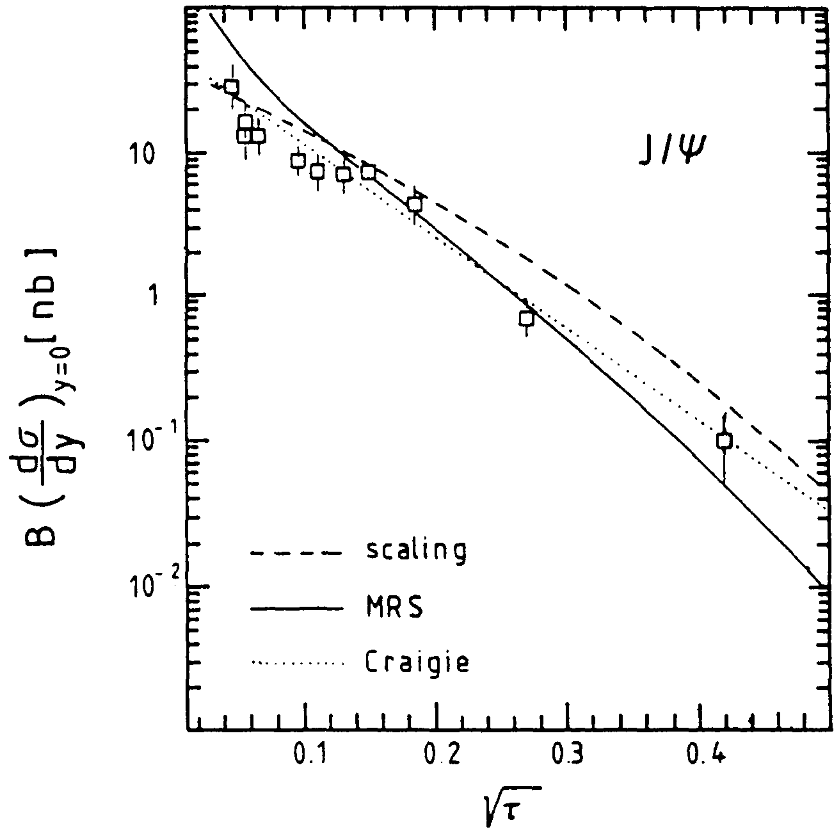

Figure 6. J/ψ production rates in $p-p$ collisions, as obtained for MRS and scaling structure functions [48] and from a fit [47], compared to experimental results [49].

second through quark-antiquark annihilation. In Figs. 5 and 6 we show $p-p$ data for these processes, compared to perturbative predictions using different structure function parametrisations. The data, available at present for CMS energies up to $\sqrt{s} \simeq 100$ GeV, are quite well accounted for. Beyond that energy, we need the structure functions at very small x. HERA will provided this; but it would certainly be good to have independent $p-p$ data at higher energies. We note further that gluon fusion as basic mechanism for charmonium production is nicely supported by recent data comparing results obtained with incident p, π^+, π^- and $\bar{p}$ beams [23].

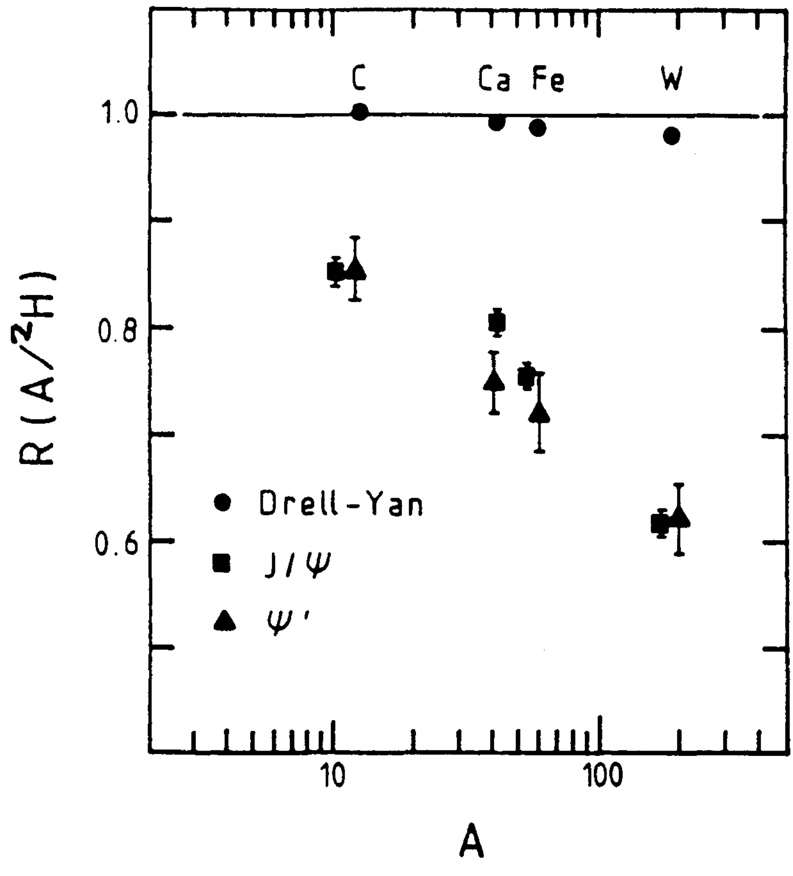

Figure 7. J/ψ and ψ' production in $p-A$ vs. $p-p$ collisions, as function of A [25].

Next we turn to hadron-nucleus collisions. It is found [24] [25] that charmonium production in such collisions (and that of bottonium as well) is suppressed with increasing A, in comparison to the corresponding production on hadrons (Figs. 7 and 8). It thus seems natural to ask if the suppression found in nucleus-nucleus collisions could be due to the same mechanism as in hadron-nucleus reactions; if this were possible, the suppression in nucleus-nucleus collisions could of course not be interpreted in terms of deconfinement. Theoretically, such a common mechanism is difficult to justify. The kinematically accessible region in $h-A$ collisions implies that the produced $c\bar{c}$ pairs leave the nucleus as a small colour octet state and reach

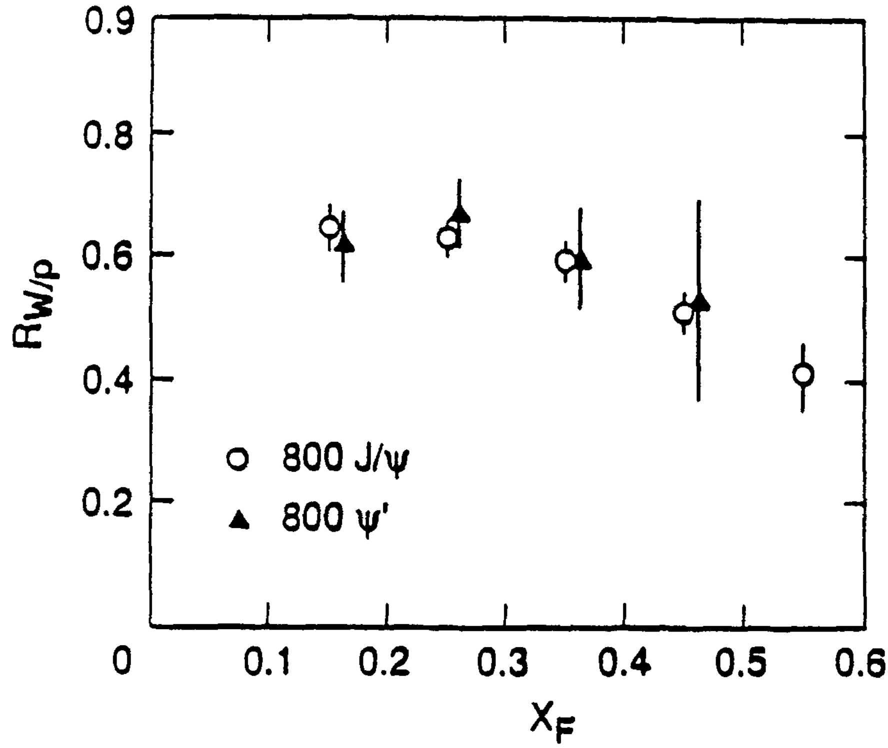

Figure 8. J/ψ and ψ' production in $p-W$ vs. $p-p$ collisions, as function of x_F [25].

physical resonance size only far outside the target; in $A-B$ collisions, on the other hand, fully formed resonances can interact with partially stopped strongly interacting matter. A specific consequence of the pre-hadronic character of the $c\bar{c}$ in $h-A$ collisions is that the suppression of the J/ψ and the ψ' should there be the same; this is in fact observed (see Figs. 7 and 8). The relevant mechanism for the suppression in $h-A$ collisions is thus expected to be a combination of colour interactions [21] [26] and gluon shadowing [27], and this does in fact account for all data in the region $0 < x_F < 1$, as seen in Figs. 9 and 10 [21].

We now return to the suppression in nucleus-nucleus collisions. As already noted, the accessible kinematic region ($0 < x_F < 0.2$) allows the formation of physical resonances in the medium. The observed suppression could thus be due to some type of absorption in confined matter at very high density [28] as well to deconfinement, and we have to find a way to distinguish between the two mechanisms. But whatever the origin of the suppression, the existence of physical resonance states predicts that the ψ' , because of its larger geometric size, will suffer a stronger suppression than the J/ψ. Specifically, this means that the ratio $\psi'/(J/\psi)$ production in $A-B$ collisions should be smaller than in $h-A$ collisions, and that it should decrease with increasing energy density (or equivalently, with E_T). Both predictions are observed [29] [30], as shown in Fig. 11. We therefore conclude that the A-dependent charmonium suppression in $h-A$ collisions and the E_T-dependent suppression in $A-A$ collisions have different origins.

The last step of our analysis should now tell us what the origin of the charmonium suppression in $A-B$ collisions is. We recall that the data [22] shows a J/ψ suppression increasing with E_T, up to a reduction of signal-above-continuum of about 50% at the highest E_T values. While the continuum was up to now mainly determined by the mass range below the J/ψ, the suppression is now established [30] also through the high mass continuum above the ψ', and it persists when the gluon shadowing effects determined in $h-A$ collisions are removed [31]. A similar but stronger suppression is found for the ψ'; moreover, the ratio $\psi'/(J/\psi)$ decreases with E_T. As is known, these features can be accounted for both by deconfinement (quark-gluon plasma formation) and by absorption in a hadronic medium of extreme density [31]. In the deconfinement picture, a charmonium state x survives if the energy density ϵ of the medium is less than the value ϵ_x needed to dissolve it. In nuclear collisions, with a hot inner and a cooler outer interaction region, the survival probability thus becomes

$$S_x^d = \left(\frac{V_{\rm cold}^x}{V}\right), \tag{1}$$

where $V_{\rm cold}^x$ is that part of the interaction volume V in which $\epsilon \leq \epsilon_x$. With ϵ proportional to the nucleon density and with standard nuclear geometry (equidistribution of nucleons in the nucleus), this gives us

$$S_x^d(\epsilon) = \Theta(\epsilon_x - \epsilon) + \left(\frac{\epsilon_x}{\epsilon}\right)^{9/4} \Theta(\epsilon - \epsilon_x) \tag{2}$$

for the survival probability of charmonium state x. In the case of absorption in an expanding medium, on the other hand, this survival probability is given by

$$S_x^a(\epsilon) = \exp\{-\sigma_x n(\epsilon)\tau_o \ln(\tau_o/\tau_f)\}; \tag{3}$$

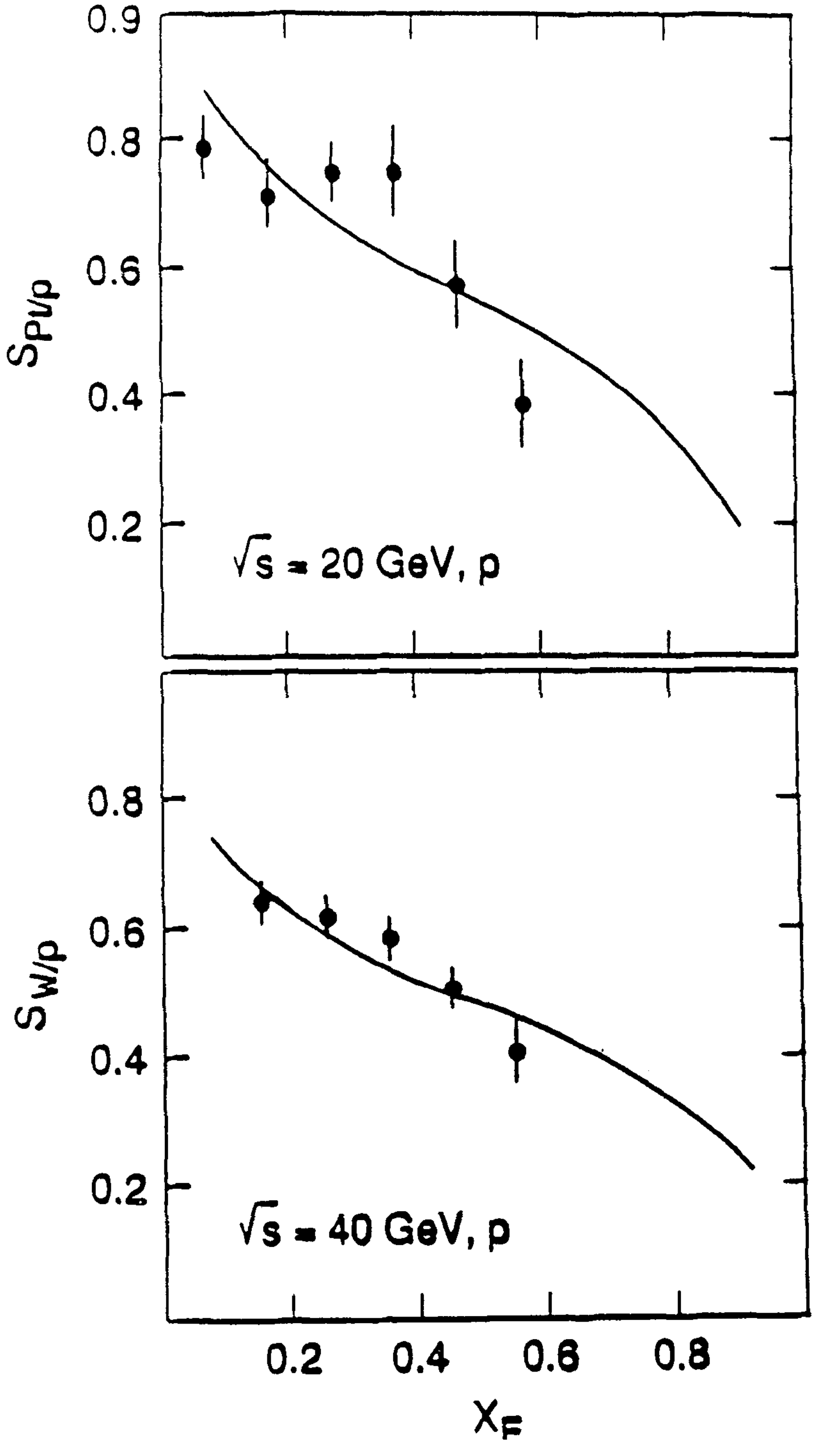

Figure 9. J/ψ production in $p - Pt$ and $p - W$ vs. $p - p$ collisions for different $\sqrt{s}$, as function of x_F [21].

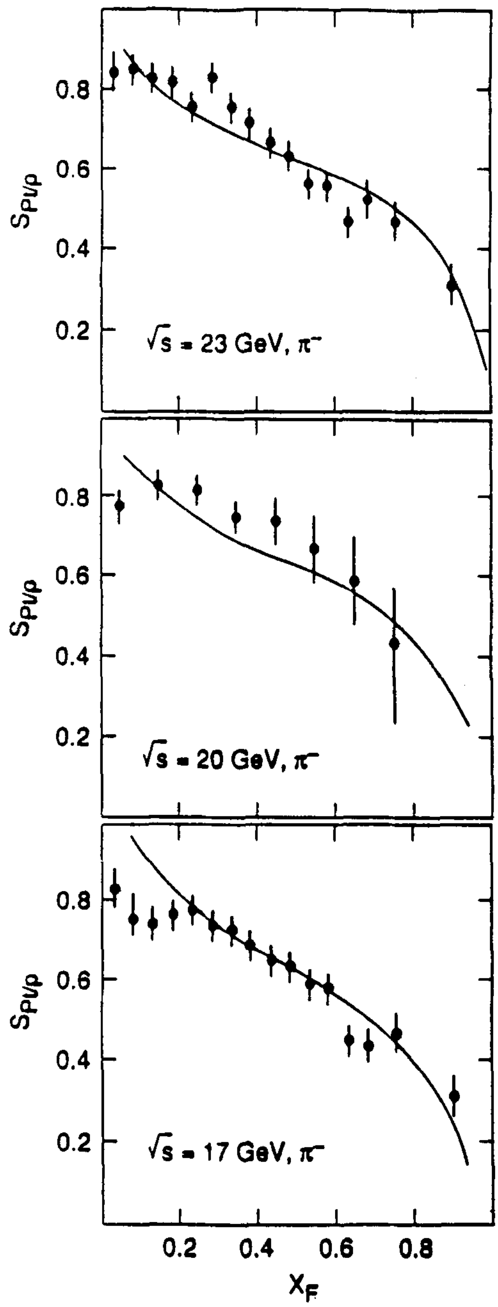

Figure 10. J/ψ production in $\pi^- - Pt$ vs. $\pi^- - p$ collisions for different $\sqrt{s}$, as function of x_F [21].

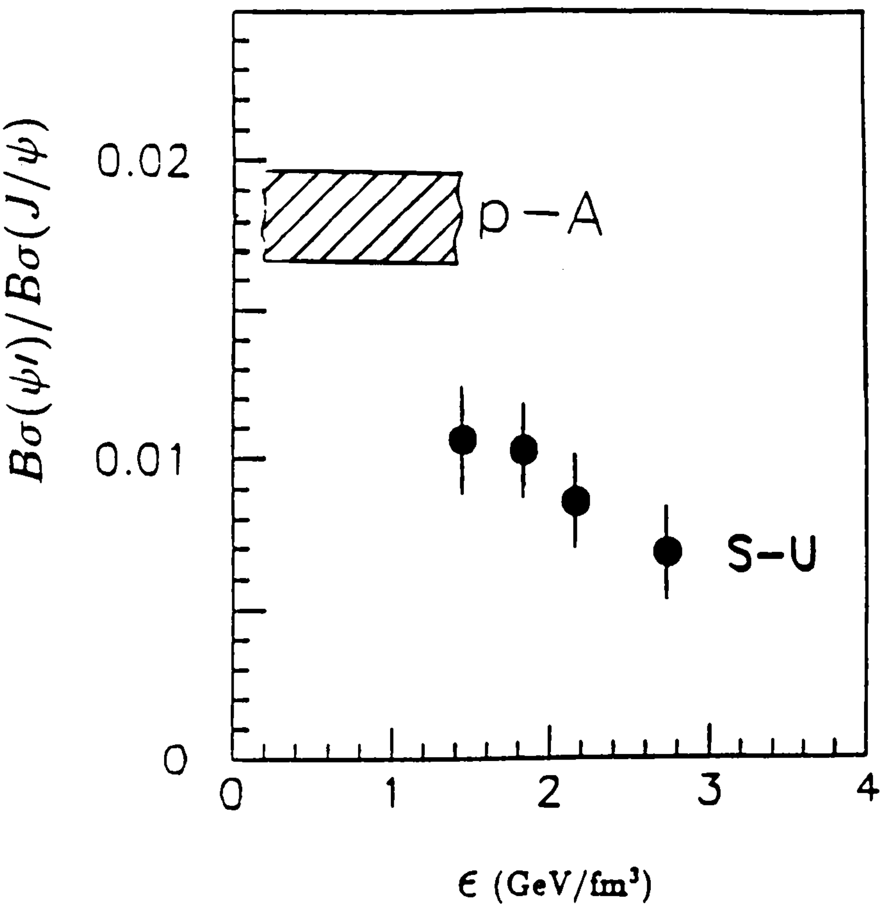

Figure 11. The ratio of ψ' to J/ψ production in $p-A$ and $S-U$ collisions as function of the initial energy density ϵ_0 [30].

here σ_x is the break-up cross section for the state x, $n(\epsilon)$ the density of absorbers at energy density ϵ, τ_o the equilibration time and τ_f the time at which the medium becomes too dilute to provide absorption.

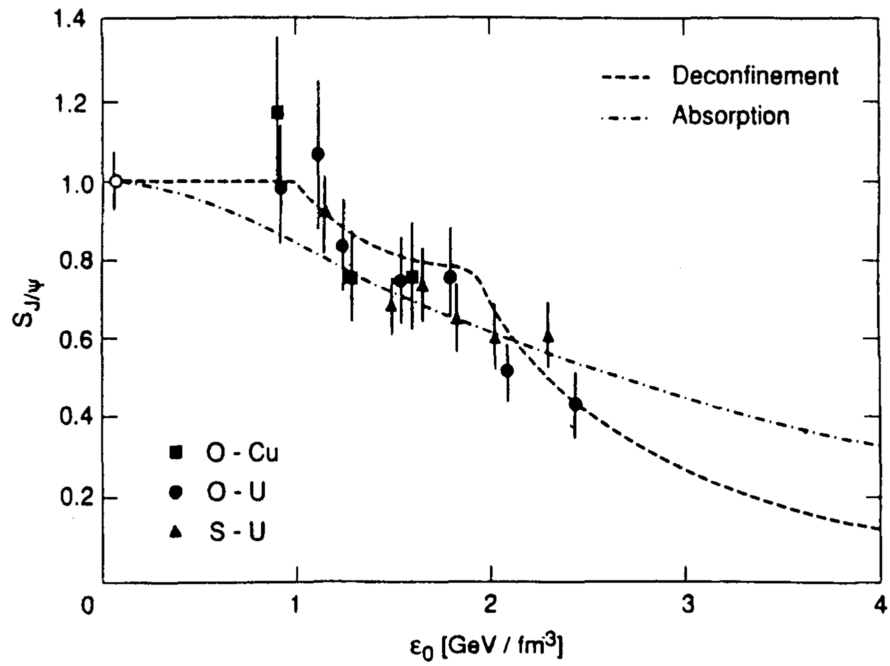

Figure 12. J/ψ suppression as function of the initial energy density ϵ_0 for deconfinement and absorption in dense matter [31].

The two J/ψ suppression patterns resulting from eqs. (2) and (3) are shown in Fig. 12, together with the presently available data [31]. Both are seen to account for the experimental results. We should note, however, that in both approaches the agreement with the data implies an energy density of 2 - 3 GeV/fm^3 at the high end. In the deconfinement picture, such a density results in a quark-gluon plasma and hence charmonium melting. In the hadronic absorption scenario, however, it means a density of four or more hadrons per normal hadronic volume [32], and the J/ψ or ψ' must interact independently with each of these hadrons. It is conceptually not clear how this can take place.

What qualitative differences are there between the two suppression mechanisms? We note two possible effects to look for.

- The pure $1S$ ψ state needs a higher energy density to "melt" than the larger excited χ_c state, which decays into a ψ. The observed J/ψ's are about 70% directly produced and about 30% due to χ-decays. The J/ψ survival probability thus becomes

$$S_J^d(\epsilon) = 0.3 \left[\Theta(\epsilon_\chi - \epsilon) + \Theta(\epsilon - \epsilon_\chi)\left(\frac{\epsilon_\chi}{\epsilon}\right)^{9/4}\right] + 0.7 \left[\Theta(\epsilon_\psi - \epsilon) + \Theta(\epsilon - \epsilon_\psi)\left(\frac{\epsilon_\psi}{\epsilon}\right)^{9/4}\right], \quad (4)$$

where the first term corresponds to the (earlier) suppression of the χ component and the second to that of the direct ψ. As a result, there is the "bump" in the deconfinement suppression pattern seen in Fig. 12, while absorption leads to a monotonic decrease in ϵ. It is quite possible that quantum effects will wash out the two-component structure with its resulting different "melting points"; but if observed, if would certainly provide evidence for deconfinement.

At any given density, larger resonances are always absorbed more than smaller; hence in the absorption scenario, the ratio $\psi'/(J/\psi)$ must vanish with increasing density. On the other hand, for a sufficiently high given density, deconfinement will melt all resonances; any survivors must come from the cool and hence non-deconfining outer regions of the interaction zone, which are determined by nuclear geometry. As a consequence, the ratio $\psi'/(J/\psi)$ takes on the constant value

$$\left(\frac{S_{\psi'}^d}{S_{J/\psi}^d}\right) = \left(\frac{\epsilon_{\psi'}}{\epsilon_\psi}\right)^{9/4} \quad \text{for } \epsilon \geq \epsilon_\psi \quad (5)$$

in the case of deconfinement, while for absorption it vanishes exponentially in ϵ, as seen in Fig. 13 [31].

At present, the statistics and the energy density range available do not yet allow any decision on the origin of charmonium suppression in $A - B$ collisions. We can conclude, however, that quarkonium production does give us a tool to probe the primordial state, and this tool has so far indicated the existence of dense primordial matter.

4. The quark-hadron transition

The first goal of high energy nuclear collisions is to produce strongly interacting *matter*. We thus have to find a way to check if the systems produced in such collisions are indeed in local thermal equilibrium. We hope that a detailed study of *thermal* hadrons (i.e., those produced in the break-up of the quark-gluon plasma) will provide a tool for this. Such hadrons will give us direct information about the system at the time they were formed: about its equilibration, size, temperature, flow structure, and more. The transition from a quark-gluon plasma to hadronic matter is one of the most fascinating features of strong interaction thermodynamics. It has attracted much interest, in statistical QCD as well as in the context of phenomenological models, and it is clearly of great importance to find ways to study it experimentally. If the quark-hadron transition and the final freeze-out

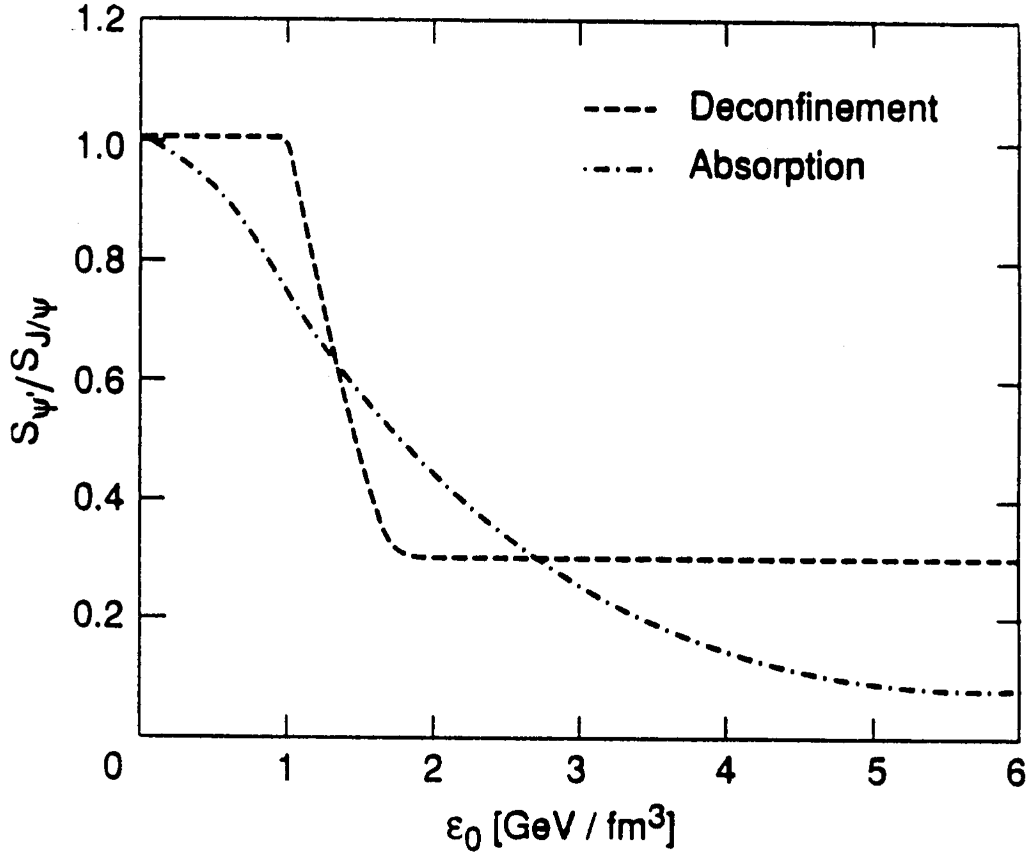

Figure 13. The ratio of ψ' to J/ψ suppression as function of the initial energy density ϵ_0 for deconfinement and absorption in dense matter [31].

are not far apart, it is quite likely that thermal hadrons will retain some transition features and thus provide us with information about the critical region. In addition, thermal hadrons can of course also give us indirect information about prior stages, if the evolution pattern of the system is known; we could then use them to determine the initial entropy and energy density or the intial temperature. Here we will not assume the evolution to be known and thus consider thermal hadron production as a way to probe the quark-hadron transition and the freeze-out era.

The basic questions here seem to be these:

- Are the finally emitted hadrons in thermal equilibrium?
- Is there one universal freeze-out point, or do strong interactions stop for some hadrons earlier than for others?

Both questions appearently deal only with hadrons made up of light quarks; charmonium and bottonium states, for example, presumably never partake in any thermal equilibrium. Because of the miniscule Boltzmann factor ($\exp-(M_J/T) \simeq 10^{-5}$ for $T \simeq 300$ MeV), the thermal production of these states is ruled out at all meaningful temperatures.

Thermal equilibrium can appear in two forms. A system containing different particle species can be in overall thermal equilibrium; this means that each species has thermal momentum distributions, and that the relative abundance of the different species is governed by their thermal weights, dependent on masses and internal degrees of freedom. Thermal equilibrium for a given species is, curiously enough, not so easy to check in high energy nuclear collisions. It implies that the distributions of the emitted mesons and baryons should be of Boltzmann form,

$$\frac{dN}{dm_T} \sim e^{-(m_T/T)}, \tag{6}$$

with $m_T = (m^2 + p_T^2)^{1/2}$ denoting the transverse energy of a hadron of mass m. It so happens that already the distributions of secondaries in $p-p$ reactions (and with some caveats, even in e^+-e^- annihilation) have this form. This either indicates that already these processes lead to thermal behaviour – a point of view explored a number of times over the past fourty years [33] – or it is due to other reasons (geometry, uncertainty effects). In either case we have no way to test thermal behaviour for a given particle species specifically for nucleus-nucleus collisions. We can, however, check if the different hadron species are produced according to some thermal weights, and this study has led to very interesting results. It started with the prediction of strangeness enhancement in nuclear collisions and was subsequently developed into a rather well-defined method to probe thermalisation in this more "chemically equilibrated" sense [2].

If at freeze-out we do indeed have an ideal gas of hadrons and hadronic resonances, then the system is described by the partition function

$$\ln Z(T,\mu_B,\mu_S) = \left[\sum_i W_i^m + \sum_i \left(l_B^{-B_i} l_S^{S_i} + l_B^{B_i} l_S^{-S_i}\right) W_i\right]. \tag{7}$$

Here W_i is the phase space factor for non-strange mesons of species i, while the second term contains all strange mesons and all baryons (both non-strange and strange), with S_i and B_i denoting the strangeness and baryon-number of the hadron in question. The baryon-number fugacity is defined as $l_B = \exp(\mu_B/T)$ and that for the strangeness as $l_S = \exp(\mu_S/T)$. The phase space factors are given by

$$W_i = \frac{d_i m_i^2 VT}{2\pi^2} K_2\left(\frac{m_i}{T}\right), \tag{8}$$

with d_i denoting the degeneracy, m_i the mass of hadron species i; V is the volume of the system. For the temperature range presently of interest ($T \leq 250$ MeV), it is sufficient to include in the sums in eq. (7) all known hadrons up to a mass of 2 GeV. The partition function (7) then determines all thermal properties of the system in terms of the three parameters T, μ_B and μ_S. One of these, say μ_S, can be eliminated by requiring the overall strangeness of the system to vanish. If there is a one-stage freeze-out of all thermal hadrons, then *all* production ratios $(K/\pi, \pi/\rho, \pi/p, K/Y, Y/\bar{Y}, \Xi/Y, ...)$ are given in terms of the two parameters T and μ_B. In other words, we can take two measured ratios, such as $\Lambda/\bar{\Lambda}$ and $\Xi/\bar{\Xi}$, to fix the values of T and μ_B, and all other ratios are then predicted, if the different particle species are indeed present according to their equilibrium weights. Although previous analyses had been in accord with this for all strange baryon and meson ratios at $T \simeq 200$ MeV [34, 38], $\mu_B \simeq 300$ MeV [39], newer and more precise data [40] no longer provide a common equilibrium freeze-out point even for strange hadrons (Fig. 14).

Bringing in non-strange hadrons had always led to disagreement, with lower freeze-out temperatures ($T \simeq 100 - 140$ MeV) as well as lower baryochemical potential ($\mu_B \simeq 200$ MeV). It is possible, however, to retain a common freeze-out for all strange hadrons by introducing partial strangeness saturation, i.e., by assuming that strangeness is not yet fully present in the system to the extent it should be if there was complete chemical equilibrium [41] With this assumption, the fugacities l_S and l_S^{-1} in eq. (7) are multiplied by a strangeness saturation factor $\gamma \leq 1$. This results in $T \simeq 200$ MeV and $\mu_B \simeq 250$ MeV as freeze-out parameters for the (partially equilibrated) strange hadrons [42], as seen in Fig. 15.

The unique freeze-out point for hyperons and cascade particles (and their antiparticles) arises by construction, of course, since we have three independent ratios and three parameters. The freeze-out point is also in accord with present Λ/K_s^o. One basic problem in this type of analysis is the availability of data in the same kinematic regions and for the same collision partners. It will be interesting to follow the further development of the freeze-out study for both strange and non-strange hadrons. All present results do suggest, however, that there is no unique freeze-out point for all hadrons.

It therefore seems reasonable to consider when freeze-out will occur. One possible definition is given by requiring the mean free path of a hadron in the medium to reach the size of the system. This leads to

$$R_F = \left[\left(\frac{3}{4\pi}\right)\left(\frac{dN}{dy}\right)\sigma\right]^{1/2} \tag{9}$$

for the freeze-out radius R_F; here σ denotes the interaction cross-section for the hadron in question. A hadron species with a bigger cross-section should thus undergo a later freeze-out [43]. If we assume strange hadrons to have an interaction cross-section half the size of non-strange cross-sections (this is the case for the Kp

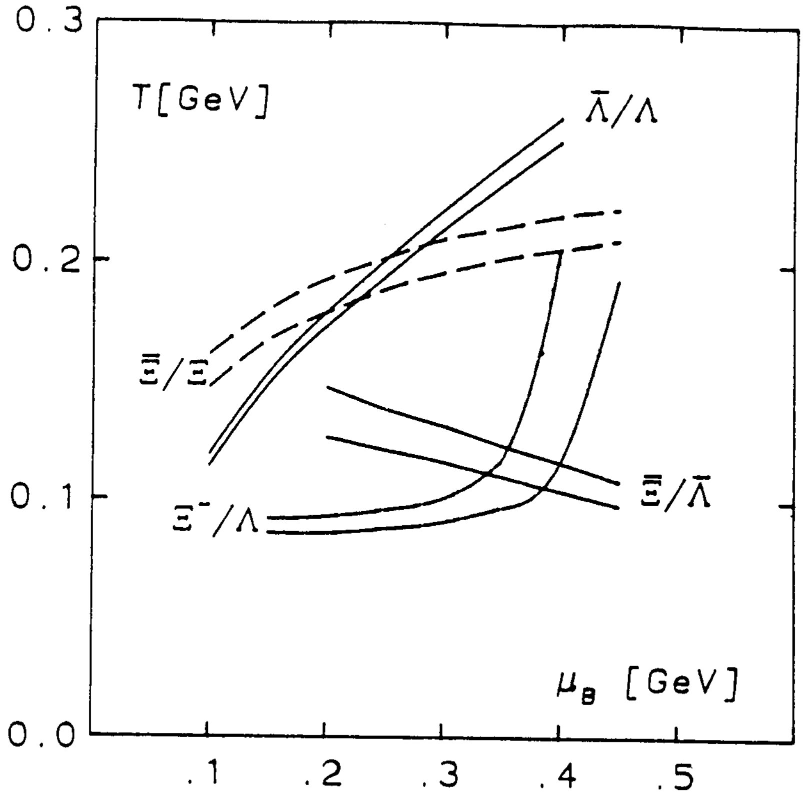

Figure 14. Strange baryon production ratios as function of temperature and baryochemical potential, with full strangeness saturation [42].

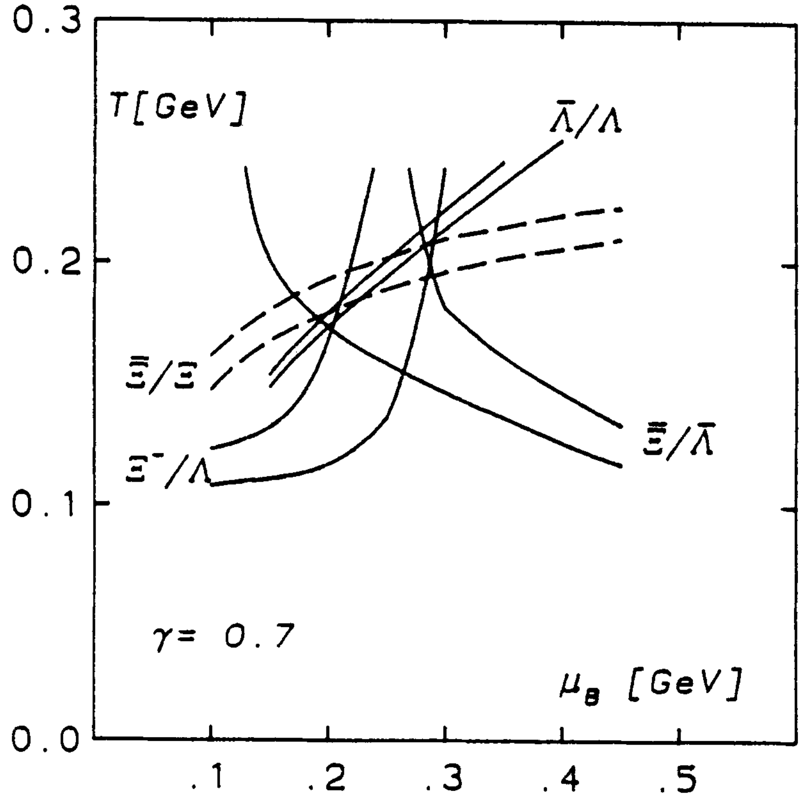

Figure 15. Strange baryon production ratios as function of temperature and baryochemical potential, with partial strangeness saturation [42].

vs. πp cross-sections at 1 GeV/c lab momentum), then we get $R_F^s \simeq 0.7\ R_F^{ns}$. This result can be checked independently by hadron interferometry studies, which determine the freeze-out radii for pions and kaons; new data [44] give at least qualitative agreement. In the case of an isentropic longitudinal expansion, these R_F-values imply $T_F^s \simeq 190$ MeV if $T_T^{ns} \simeq 130$ MeV, in reasonable accord with what we found above. Hadron spectra thus seem to suggest that freeze-out is a sequential process, determined by the different hadronic interaction cross-sections; however, this conclusion is at present still of a very qualitative nature.

Nevertheless, we have seen that hadron spectra and correlations provide us with a good probe for the study of hadronisation and the quark-hadron transition region.

5. Conclusions

We have seen that different signals probe different stages in the evolution of the quark-gluon plasma. It is therefore necessary to clarify what we can learn from any given signal, and from what stage it brings us information. We saw that high mass Drell-Yan pairs and open charm mesons probe the pre-equilibrium stage; the suppression of charmonium and bottonium production provides a tool for the study of the hot primordial QGP; hadronic spectra and correlations give us information about the quark-hadron transition and freeze-out. Clearly they are just some of the possible signals and probes; there are many others – thermal dileptons, direct photons, jet production, hadron mass modifications in media, flow effects, strangelets, to name some that have so far been studied quite extensively. We will certainly need all of them in order to establish the complete evolution pattern of the quark-gluon plasma.

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