# J/ψ SUPPRESSION BY QUARK–GLUON PLASMA FORMATION ☆

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If high energy heavy ion collisions lead to the formation of a hot quark–gluon plasma, then colour screening prevents $c\bar{c}$ binding in the deconfined interior of the interaction region. To study this effect, the temperature dependence of the screening radius, as obtained from lattice QCD, is compared with the J/ψ radius calculated in charmonium models. The feasibility to detect this effect clearly in the dilepton mass spectrum is examined. It is concluded that J/ψ suppression in nuclear collisions should provide an unambiguous signature of quark–gluon plasma formation.

Statistical QCD predicts that strongly interacting matter should at sufficiently high density undergo a transition from hadronic matter to quark–gluon plasma [‡1]. It is hoped that energetic nuclear collisions will allow us to study this transition in the laboratory [‡2]. The experimental detection of plasma formation thus becomes crucial: what observable signatures does the predicted new form of matter provide?

Signatures proposed so far include [‡3] real or virtual photons, the $p_T$ distribution of secondary hadrons, and the relative production rate of strange particles. Non-thermal processes as well as uncertainties in the plasma evolution do, however, lead to considerable ambiguity for the signals considered up to now. We want to present here another type of signature for plasma formation, which directly reflects deconfinement and appears to provide a rather clear and model-independent test.

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‡1 For a recent survey see ref. [1].

‡2 For a recent survey see ref. [2].

‡3 For surveys see ref. [3].

The basic mechanism for deconfinement in dense matter is the Debye screening of the quark colour charge [4]. When the screening radius $r_D$ becomes less than the binding radius $r_H$ of the quark system, i.e., less than the hadron radius, the confining force can no longer hold the quarks together and hence deconfinement sets in. We shall investigate here the effect of such a deconfining medium on the binding of c and $\bar{c}$ quarks into J/ψ mesons.

The temperature dependence of the colour screening radius was recently studied in SU(2) [5] and SU(3) [6] gauge theory. There, one considers the interaction of a static quark–antiquark system in a purely gluonic thermal environment. The absence of dynamical quarks does, of course, change the screening phenomenon considerably [5]: since the quarks transform according to the fundamental representation of the colour gauge group and the gluons according to the adjoint, the quark colour charge cannot be screened directly. Nevertheless, the quark interaction is mediated by gluons, and at high temperature the dominant contribution will come from the exchange of one gluon, made massive by gluonic colour screening. Moreover, we expect that the introduction of dynamical quarks will, if anything, enhance the screening, as it increases the density of colour-carrying constituents.

The quark–antiquark interaction in SU($N$) gauge theory is parameterized by the correlation function $\Gamma(r, T)$, where $r$ denotes the distance of separation for the static $q\bar{q}$ system and $T$ the temperature of the gluonic heat bath. For large $r$, $\Gamma(r, T)$ decreases exponentially,

$$\Gamma(r, T) \sim \exp[-r/\xi(T)] , \qquad (1)$$

with $\xi(T)$ denoting the correlation length. Its temperature dependence is shown in fig. 1, as obtained from the correlation of Polyakov loops in lattice gauge theory [5,6]. To convert the lattice results into physical units, we have fixed the deconfinement temperature $T_c = 200$ MeV. We note that $\xi$ drops quite rapidly with $T$ and that at $T/T_c = 1.5$, is of the order of 0.2–0.3 fm. From what was said above, we expect

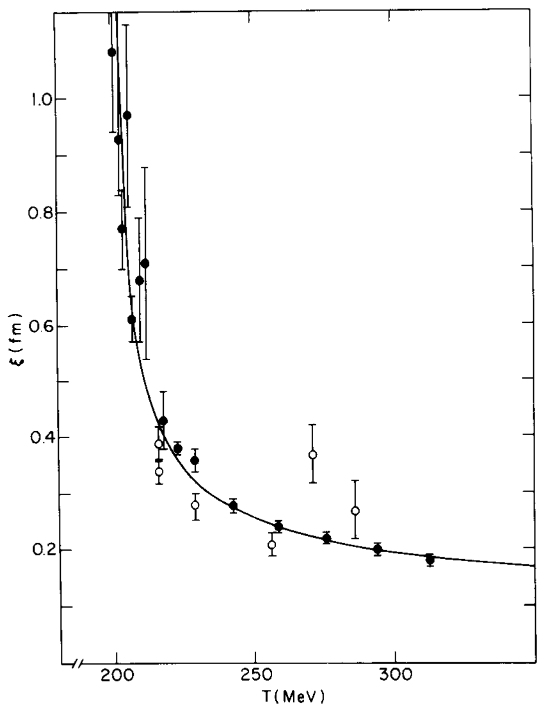


Fig. 1. Temperature dependence of the correlation length, as obtained in SU(2) gauge theory (solid dots, from ref. [5]) and in SU(3) gauge theory (open circles, from ref. [6]); here $T_c = 200$ MeV was used to fix the scale.

this to be an upper bound for the colour screening radius $r_D(T)$ in QCD with dynamical quarks. In particular, in full QCD, there will be colour screening between quarks even at $T_c$, whereas the static quarks in pure gauge theory experience at that point an effectively unscreened three-dimensional Coulomb field.

Let us now consider $J/\psi$ production, first in hadron–hadron collisions. The dominant mechanism [7] is hard parton–parton interaction, producing $c\bar{c}$ pairs. The subsequent resonant interaction of the $c\bar{c}$ system then leads to $J/\psi$ production. If, however, the $c\bar{c}$ production occurs in a nuclear collision, and if such collisions result in a quark–gluon plasma, then the produced $c\bar{c}$ finds itself in a deconfining environment. Provided the temperature of this environment is sufficiently high – i.e., provided the screening radius $r_D(T)$ is smaller than the binding radius $r_{J/\psi}(T)$ – then the resonance interaction cannot become operative and $J/\psi$ production will be prohibited. The c and the $\bar{c}$ will proceed on separate trajectories aand eventually lead to the production of "open charm" mesons ($c\bar{u}$, $\bar{c}u$, etc.).

To assure that this $J/\psi$ suppression in nuclear collisions indeed constitutes an observable signature of plasma formation, we must answer a number of questions:

(i) Can the $J/\psi$ escape from the production region before plasma formation?

(ii) At what temperature does $r_D(T)$ fall below $r_{J/\psi}(T)$, and how does $r_{J/\psi}(T)$ behave as function of $T$? The large mass gives the $J/\psi$ a smaller radius than that of conventional mesons, and sufficiently small hadrons could survive deconfinement as coulombic bound states until much higher temperatures.

(iii) Are there competitive non-plasma $J/\psi$ suppression mechanisms?

(iv) Could the $J/\psi$ suppression in the plasma be compensated in the transition or hadronization stage?

(v) Could enhanced non-resonant production of lepton pairs ("thermal dileptons") prevent the observation of the $J/\psi$? In this case, we could not study deconfinement directly, although plasma formation would still be the cause for not seeing $J/\psi$'s.

We will now take up these questions.

The time $\tau_0$ after the collision, which is necessary to form a plasma in something like thermal equilibrium, is expected to be of the order of one fermi [8].

Certainly the hard production of $c\bar{c}$ pairs occurs at times $\tau \ll \tau_0$. However, to form a $J/\psi$ with its intrinsic dimension from this $c\bar{c}$ pair will again require a time of order $\tau_0$, unless $r_{J/\psi}(T)$ is very much smaller than the typical hadronic scale. We shall see shortly that this is not the case. Hence the $J/\psi$'s cannot appear before plasma formation. In addition, anything produced in the interior of a nuclear interaction region still has to travel at least a distance of about $A^{1/3}$ fm before it could get out; here $A$ denotes the nuclear mass number.

Next, we want to look at the radius of the $J/\psi$. Charmonium models [9] suggest for the $c\bar{c}$ system a non-relativistic interaction potential

$$V(r) = \sigma r - \alpha_{\text{eff}} r\,, \tag{2}$$

where $\sigma$ is the string tension and $\alpha_{\text{eff}}$ the coulombic interaction coupling. For an isolated $c\bar{c}$ system (i.e., at $T=0$), typical values are $\sigma \simeq 0.16\ \text{GeV}^2$, $\alpha_{\text{eff}} \simeq 1/2$. The energy of the bound state may be estimated semi-quantitatively by

$$E(r) = 2m + 1/2mr^2 + V(r) \tag{3}$$

including the c-quark rest masses $m$ and their kinetic energy. To find the lowest state, we minimize $E(r)$ and obtain

$$1/mr^3_{J/\psi} - \alpha_{\text{eff}}/r^2_{J/\psi} - \sigma = 0 \tag{4}$$

as relation between the $J/\psi$ radius $r_{J/\psi}$ and the parameters $m, \sigma$, and $\alpha_{\text{eff}}$. With the values for $\sigma$ and $\alpha_{\text{eff}}$ quoted above, eq. (3) has a minimum at $E=3.1$ GeV, if we set $m=1.56$ GeV; this gives $r_{J/\psi} \simeq 0.20$ fm. Other reasonable parameter values (smaller $\alpha_{\text{eff}}$, slightly smaller $m$) tend to produce a somewhat larger $r_{J/\psi}$; $0.2 \leqslant r_{J/\psi} \leqslant 0.5$ fm is generally considered to be the typical range. These values all agree quite well with what would be obtained for $\alpha_{\text{eff}}=0$; this leads to $r_{J/\psi} = (m\sigma)^{-1/3}$. At $T=0$, the $J/\psi$ radius is thus largely determined by the confining part of the potential; although somewhat smaller than the radius of conventional mesons, $J/\psi$ is still of hadronic size.

With increasing temperature, $\sigma(T)$ decreases, and at deconfinement $\sigma(T_c)=0$. For $T \geqslant T_c$ we thus expect [5]

$$V(r) = -(\alpha_{\text{eff}}/r)\exp[-r/r_D(T)] \tag{5}$$

as colour-screened coulombic potential. This potential could, however, still provide bound states. Inserting the form (5) into eq. (3) and minimizing $E(r)$, we get

$$x(x+1)\exp(-x) = (m\alpha_{\text{eff}} r_D)^{-1} \tag{6}$$

with $x \equiv r/r_D$, as condition for a bound state. Eq. (6) has a solution only if $(m\alpha_{\text{eff}} r_D)^{-1} \leqslant 0.84$, so that

$$r_D^{\min} = [0.84\, m\alpha_{\text{eff}}(T)]^{-1} \tag{7}$$

is the smallest value of the screening radius still permitting a coulombic bound state. With the $T=0$ value $\alpha_{\text{eff}}=1/2$ we get $r_D^{\min} \simeq 0.31$ fm; if we take into account the temperature decrease of $\alpha_{\text{eff}}$, we obtain considerably larger values. The lattice calculations of ref. [7] give at $T/T_c=1.5$ the value $\alpha_{\text{eff}} \simeq 0.2$; then bound states would already become impossible at this temperature for $r_D^{\min} \simeq 0.76$ fm.

From eq. (6) we get

$$(r^{\max}_{J/\psi}/r_D) = 1.61 \tag{8}$$

as the universal coulombic $J/\psi$ radius at the last point where such a state still exists. For the $\alpha_{\text{eff}}$ values considered above, this implies values in the range $0.5 \leqslant r^{\max}_{J/\psi} \leqslant 1.3$ fm for the size of the coulombic $c\bar{c}$ bound state just before it disappears. The $J/\psi$ has thus already become quite large; the shift in overall mass, however, remains quite small, due to the heavy c-quarks.

Comparing our results with behaviour of $\xi(T)$ in fig. 1, we conclude that at $T/T_c=1.5$, where $\xi \simeq 0.2$ fm, the production of $J/\psi$'s is not possible, even as coulombic bound state. Taking the limiting value $r_D^{\min} \simeq 0.31$ fm at face value, the existence of $J/\psi$'s is excluded even down to $T/T_c=1.2$ or less. Plasma formation thus prevents $J/\psi$ formation already just above $T_c$.

We note here that the effect of decreasing confinement on the $c\bar{c}$ binding of the $J/\psi$ below the $T_c$ was recently studied [10] and shown to provide a shift to lower $J/\psi$ mass, The crucial point of our result is that for $T$ just above $T_c$, the $J/\psi$ will completely disappear in the deconfining plasma. It does not imply, as we shall see, that thermal emission is the dominant source of dileptons in the $J/\psi$ region. A situation of hard interactions dominating the background dilepton continuum in the $J/\psi$ region still remains and thus makes the direct observation of $J/\psi$ deconfinement feasible.

We have concentrated on the $J/\psi$, as the most striking resonance signal observed in the lepton pair spectrum. Since the $\psi'$ radius is presumably slightly larger than that of the $J/\psi$, its production should of course be suppressed as well.

Next we will address the question of alternative suppression mechanisms. It is possible that not only plasma formation, but also some type of nuclear absorption would prevent the $J/\psi$ signal from appearing in nuclear collisions? Here we recall the experimental studies of $J/\psi$ production in photoproduction [11] and proton–nucleus [12] processes: they indicate that there is essentially no nuclear absorption of the $J/\psi$. Correspondingly, the $J/\psi$–nucleon cross section is only about 1–3 mb, compared to the 40 mb for pp interactions. Incidentaly, this much reduced strong interaction of the $J/\psi$ has led to some models [13] proposing the use of $J/\psi$'s as primordial plasma signal. Our considerations, on the contrary, exclude primordial $J/\psi$ formation in a deconfining medium.

The following two questions – $J/\psi$ production during and after the hadronization transition, and thermal dileptons – are rather related. There are three distinct types of dilepton production: the hard quark–antiquark annihilation of the Drell–Yan mechanism, thermal dilepton production in a hot quark–gluon plasma, and hadronic dilepton formation, e.g., through vector meson dominance. In hadron–hadron interactions, the first of thse mechanisms provides dileptons in the $J/\psi$ mass region and above, while the last dominates for low mass pairs in the $\rho$, $\omega$, $\phi$ region. Thermal dilepton emission is already a plasma process and has been proposed as "thermometer" for the plasma temperature [14]. Its observation would in itself provide an indication of plasma formation; if abundantly produced in the $J/\psi$ region, it could, however, mask the deconfinement phenomenon we wish to study here. Therefore we would like to find a kinematic region which allows plasma formation, yet provides in the $J/\psi$ mass region dominantly Drell–Yan dilepton production. In this case, the $J/\psi$ could and would be formed, unless the deconfining plasma prevents $c\bar{c}$ binding.

The dilepton spectrum from the Drell–Yan mechanism has the form

$$(\mathrm{d}^2\sigma/\mathrm{d}M^2\,\mathrm{d}y)^{\mathrm{DY}}_{\mathrm{pp}} = (1/M^4)f(M/\sqrt{s}) \quad [\mathrm{mb\ GeV^{-2}}]\,, \tag{9}$$

here $\sqrt{s}$ is the CMS collision energy and $M$ the lepton pair mass. Proton–proton and proton–nucleus data, the latter scaled by $A^{-1}$, give for $M \gtrsim 3$ GeV in the central region ($y \sim 0$) the form [7]

$$f_{\mathrm{pp}}(M/\sqrt{s}) \simeq (3.75\times10^{-5}) \times \exp(-15\,M/\sqrt{s})\,. \tag{10}$$

The average lepton pair yield per pp collision is given by

$$(\mathrm{d}^2N/\mathrm{d}M^2\,\mathrm{d}y)^{\mathrm{DY}}_{\mathrm{pp}} = (1/\sigma^{\mathrm{pp}}_{\mathrm{in}})(\mathrm{d}^2\sigma/\mathrm{d}M^2\,\mathrm{d}y)^{\mathrm{DY}}_{\mathrm{pp}}\,, \tag{11}$$

where $\sigma^{\mathrm{pp}}_{\mathrm{in}}$ is the elastic pp cross section. The total numbers of Drell–Yan lepton pairs emitted in a central collision of nuclei $A$ and $B(A>B)$ may be estimated by multiplying (11) by the total number of effective nucleon–nucleon collisions $\nu(A,B)$,

$$\nu(A,B) \simeq (2R_{\mathrm{A}}/\lambda_{\mathrm{in}})B = 2R_{\mathrm{A}}\rho_0\sigma^{\mathrm{pp}}_{\mathrm{in}}B\,, \tag{12}$$

where $\lambda_{\mathrm{in}} = (\sigma^{\mathrm{pp}}_{\mathrm{in}}\rho_0)^{-1}$ is the inelastic mean free path of a nucleon traversing nuclear matter at density $\rho_0 = 0.16\ \mathrm{fm}^{-3}$. For the present semi-quantitative analysis, we use the estimate

$$(\mathrm{d}^2N/\mathrm{d}M^2\,\mathrm{d}y)^{\mathrm{DY}}_{\mathrm{AB}} \simeq (A^{1/3}B/M^4)(3.6\times10^{-7}) \times \exp(-15M/\sqrt{s}) \quad [\mathrm{GeV^{-2}}]\,, \tag{3}$$

obtained by multiplying eq. (11) with $\frac{3}{4}\nu$, to account for geometric effects, which reduce the Drell–Yan yield when $A \approx B$.

By using the experimental form (10) to determine this rate, we have automatically taken into account the correction which would otherwise be needed [15] to fit the pp data to Drell–Yan calculations based on deep inelastic structure function measurements.

The thermal dilepton spectrum from a hot quark–gluon plasma undergoing longitudinal scaling hydrodynamic expansion [16] is given by

$$\mathrm{d}^2N/\mathrm{d}M^2\,\mathrm{d}y)^{\mathrm{TH}}_{\mathrm{AB}} \simeq B^{2/3}(25\alpha^2/\pi^2)\sqrt{\tfrac{1}{2}\pi}\,(T_0/M)^{1/2} \times \exp(-M/T_0) \quad [\mathrm{GeV^{-2}}]\,, \tag{14}$$

where $T_0$ denotes the initial temperature of the plasma at the time of its formation; here $\alpha = \frac{1}{137}$. We have in eq. (14) assumed $M \gg T_0 > T_{\mathrm{c}}$. The form (14) is obtained from that derived in ref. [17] by use of the relation

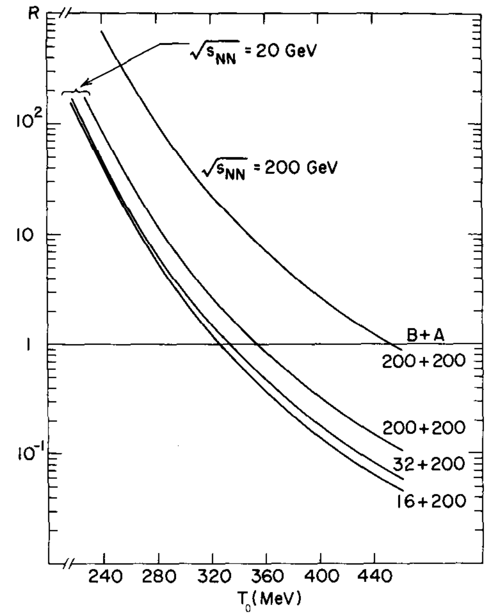


Fig. 2. Ratio of Drell–Yan to thermal lepton pair production at pair mass $M=3.1$ GeV, as function of initial plasma temperature $T_0$, for different incident nuclei and energies.

$$\tau_0 T_0 \simeq 1 \tag{15}$$

between $T_0$ and the plasma formation time $\tau_0$.

As mentioned, we want to find a regime in which there is plasma formation, but yet still Drell–Yan dominance in the $J/\psi$ region. Let us therefore compare the Drell–Yan rate with that for thermal dileptons. Here, as before, we consider the central region $y \sim 0$. From eqs. (13) and (14) we get the ratios

$$R \equiv (d^2N/dM^2dy)^{DY}_{AB}/(d^2N/dM^2dy)^{TH}_{AB}$$
$$\simeq [(2.13\times10^{-3})(AB)^{1/3}/M^4](M/T_0)^{1/2}$$
$$\times \exp[M(1/T_0 - 15/\sqrt{s})] . \tag{16}$$

In fig. 2. we show $R$ as a function of $T_0$ at $M=3.1$ GeV for several values of $A$ and $B$. We note that even for CERN-SPS experiments with $^{16}O$ and $^{32}S$ beams, there exists a "window" in which the Drell–Yan mechanism dominates thermal emission. This dominance is greatly enhanced for the planned relativistic heavy ion collider (RHIC) at BNL, with colliding $A=200$ beams at $\sqrt{s_{NN}}=200$ GeV. Incidentally, a further increase of $s_{NN}$ above this value would not change $R$ much more. In general, $R$ increases with $A$, $B$ and (up to saturation) with $s$; it decreases with $T_0$.

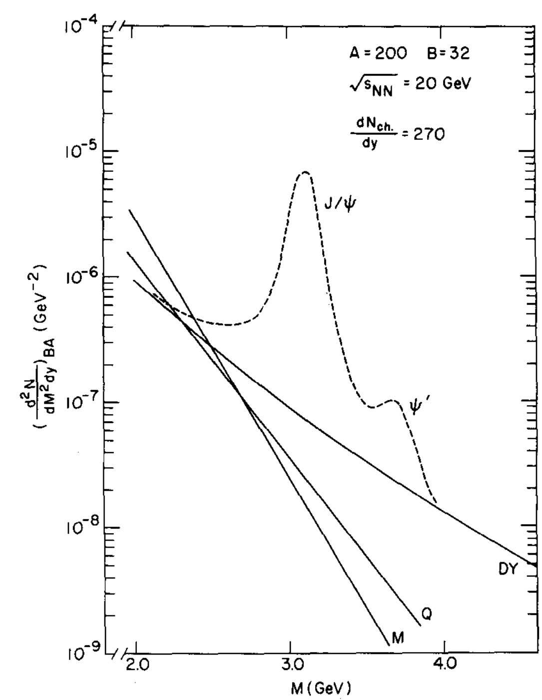


Fig. 3. Lepton pair production by Drell–Yan mechanisms (DY), thermal emission from a plasma at $T_0=300$ MeV (Q) and thermal emission from the transition region at $T_c=200$ MeV (M). The dashed line indicates the unsuppressed, scaled-up $J/\psi$ and $\psi'$ signal. The collision parameters are chosen for the planned CERN-SPS experiment.

To illustrate the behaviour of the spectrum in the $J/\psi$ region, we show in figs. 3 and 4 the forms (13) and (14) for the $^{32}S$ beam at the SPS and for the RHIC parameters, both at $T_0=300$ MeV. Included here is also the result for thermal emission from the transition region [18]; in the $J/\psi$ mass range at $T_0=300$ MeV, it falls below the plasma yield. It is evident from figs. 3 and 4 that the Drell–Yan dominance increases with $M$. Hence to assure that this mechanism prevails in a given experiment, it seems best to first measure the spectrum at large $M$. If its extrapolation into the $J/\psi$ region agrees with the data there, then we can assume that it dominates there as well. In this case, the suppression of the $J/\psi$ would provide a clear test for deconfinement.

It should be noted that although plasma formation forbids $J/\psi$ production in the interior of the interac-

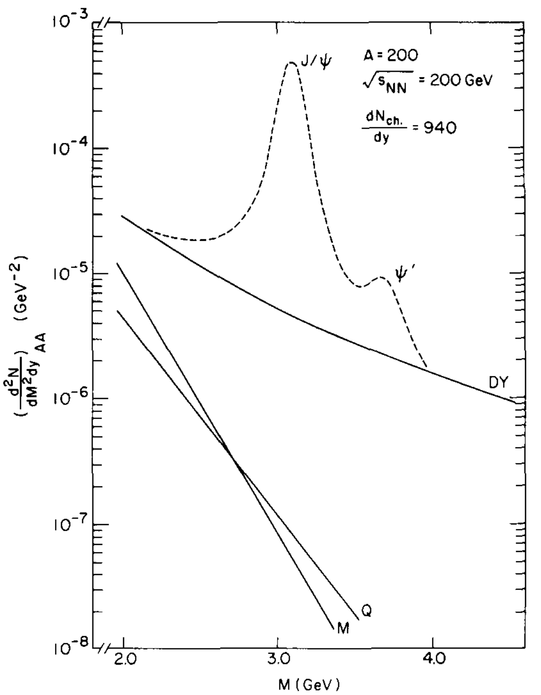


Fig. 4. Same as fig. 3, but for the collision parameters of the planned RHIC-facility of BNL.

tion region, it remains possible at the transverse perimeter of the nuclei, where we expect essentially nucleon–nucleon collisions. Hence we expect strong but not total suppression for actual $AB$ collisions – $J/\psi$ production can still occur, but at a rate decreasing as $B^{-2/3}$ (with $A > B$). For light ions, such as $^{16}O$, the suppression may thus be considerably weaker. The same holds true for peripheral interactions – the suppression is highest for head-on collisions.

In the above figures, $T_0$ was treated as a given parameter. Let us recall its relation to the observable central charged-particle multiplicity ($dN_{ch}/dy$). For isentropic expansion we have [17,19]

$$dN_{ch}/dy = \tfrac{3}{4}\cdot\tfrac{1}{4}\, dS/dy = \tfrac{3}{4}\cdot\tfrac{1}{4}\, \pi R^2 \tau_0 (\tfrac{4}{3}\cdot 15.6\, T_0^3)$$

$$\simeq 12(T_0/200\ \mathrm{MeV})^2 B^{2/3}\,, \qquad (17)$$

so that $T_0$ increases as the square root of the charged-particle multiplicity. We have here again taken $\tau_0 T_0 = 1$. To obtain ($dN_{ch}/dy$), we have multiplied the overall $\pi$–K multiplicity by $\frac{3}{4}$; the factor 15.6 counts three equivalent quark flavours (u, d, s). The resulting charged-particle multiplicities are indicated in figs. 3 and 4. If we scale them by $B$ to obtain an indication of the coresponding multiplicities in pp interactions, we obtain 8–9 and 4–5 for figs. 3 and 4, respectively. The measurement of the charged multiplicities implied by the parameters of figs. 3 and 4 should therefore be possible.

Our argument for $J/\psi$ suppression as plasma signature consists of two parts. The absence of $J/\psi$'s in a deconfined medium above some determinable temperature is a general and model-independent phenomenon. To check if this phenomenon is observable in nuclear collisions, we must at present make use of models for such processes. We have done so not because we wanted to fix the precise range for the observation of this phenomenon; nuclear collision theory is not yet sufficiently quantitative for that. Instead, we wanted to indicate that for what is considered a reasonable description of heavy ion collisions there is indeed a range where $J/\psi$ suppression should become evident. The final parameters of this range will very likely have to be fixed on the basis of more detailed experimental information (Drell–Yan region, collision centrality, etc.).

We thus conclude, that there appears to be no mechanism for $J/\psi$ suppression in nuclear collisions except the formation of a deconfining plasma, and if such a plasma is produced, there seems to be no way to avoid $J/\psi$ suppression. Furthermore, our estimates indicate that the measurement of the dilepton spectrum from nuclear collisions should allow a clear test of this phenomenon.

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