TRANSVERSE ENERGY DEPENDENCE OF J/ψ SUPPRESSION IN HEAVY ION COLLISIONS

F. KARSCH
Theory Division, CERN, CH-1211 Geneva 23, Switzerland

and

R. PETRONZIO
Department of Physics, University of Rome II, Tor Vergata, via Orazio Raimondo, I-00173 Rome, Italy

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We discuss the transverse energy and the nuclear size dependence of J/ψ suppression in heavy ion collisions under the assumption that a quark–gluon plasma has been formed.

1. Introduction

There is experimental evidence for a suppression of J/ψ production in heavy ion collisions associated with large transverse energy events [1]. The amount of suppression strongly depends on the transverse momentum p_T of the J/ψ and the total transverse energy E_T of the event. These effects are expected to occur as a consequence of quark–gluon plasma formation [2–4]. Given our ignorance of complicated nuclear effects, the interpretation in terms of a quark–gluon plasma cannot be claimed to be unique, but it is simple.

Attempts to understand J/ψ suppression in terms of more conventional nuclear absorption or hadronic gas models [#1] have shown that one can get a large amount of suppression also in these models. However, in general these approaches lead to a p_T dependence flatter than in the plasma hypothesis. At present the experimental data are not accurate enough to distinguish between both mechanisms and it is thus important to consider further observables that may allow to discriminate between different approaches.

In a recent paper we have analyzed the momentum dependence of J/ψ suppression under the assumption that this effect is due to formation of a quark–gluon plasma [3,4]. Although there enter many parameters in such an analysis which are not well known, we found that the observed momentum dependence of the suppression pattern is well reproduced with a reasonable choice of these parameters. The suppression pattern will, however, change with transverse energy. In the plasma model this corresponds to a change of the initial plasma temperature which in turn controls the plasma lifetime and thus the efficiency of the plasma as a screening medium.

All parameters of the plasma model can be arranged by analyzing the suppression pattern in a given transverse energy bin. This fixes completely the transverse energy dependence of the suppression pattern which will be discussed in this letter. We will use here all the parameters for the plasma model discussed in ref. [4]. In the next section we derive an expression for the total energy density in the plasma phase at some initial time t_i and relate it to the plasma lifetime t_f using a simple hydrodynamical evolution model for the plasma phase. In section 3 we use this to determine the survival probability of J/ψ resonances as a function of E_T. We obtain a prediction for the p_T dependence of the suppression pattern in different E_T-bins.

[#1] For a complete list of references see ref. [5].

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2. Energy density in the plasma phase

Our basic assumption is that in a sufficiently violent nucleus–nucleus collision a quark–gluon plasma forms which reaches equilibrium after a time $t_i \simeq 1$ fm. At the initial time t_i the plasma covers a cylindrical region around the collision axis. Its radial extension is some fraction of the radius of the projectile nucleus. Due to the non-uniform density distribution in the incident projectile nucleus as well as in the target nucleus also the energy density distribution will not be uniform. Like in our analysis of the p_T dependence of J/ψ suppression we assume that this can be parametrized by a temperature profile function of the form [4]

$$T(r) = T(0)[1-(r/R)^2]^{b/3}\,, \tag{1}$$

with $R = 1.2A^{1/3}$ being the projectile nucleus radius. Depending on the relative size of the target and projectile nuclei the parameter b may vary between 1/4 and 1/2. In the following we will use $b = 1/3$ [4].

Eq. (1) describes the radial dependence of the temperature at some initial time $t_i \simeq 1$ fm at which an equilibrated plasma may have been formed in the central region. This determines also the radial extension of the plasma region: it covers the central region up to the point where the temperature drops below the transition temperature T_c, i.e. up to

$$R_c = R\sqrt{1-(T_c/T_i)^{3/b}}\,, \tag{2}$$

where T_i denotes the initial temperature at time t_i and at the origin $r = 0$.

Lattice simulations suggest that the energy density in a quark–gluon plasma is well approximated by that of an ideal quark–gluon gas even for temperatures close to T_c [2]. For the energy density at the initial time t_i we then assume

$$\epsilon(r) = \tfrac{1}{15}\alpha\pi^2 T^4(r)\,, \tag{3}$$

with α counting the effective number of degrees of freedom in the plasma phase, i.e.

$$\alpha = N^2 - 1 + \tfrac{7}{4}Nn_f = 18.5 \tag{4}$$

for SU(N), $N = 3$, and $n_f = 2$ light quarks [3].

A rather uncertain parameter is the longitudinal extension l of the plasma region and we leave it for the moment as a free parameter. The plasma covers a cylindrical volume of size $\pi R_c^2 l$ at the initial time t_i. We thus obtain

$$E = \tfrac{1}{15}\pi^3[3\alpha/(4b+3)]lR^2T_i^4$$
$$\times[1-(T_c/T_i)^{(4b+3)/b}]\,, \tag{5}$$

for the total energy in the plasma phase. We note that E is proportional to $A^{2/3}$ and T_i^4. The initial temperature T_i fixes the lifetime of the plasma phase and this in turn determines to a large extent the suppression rate of heavy resonances [4].

Assuming longitudinal hydrodynamic expansion of the plasma [8] we can relate the plasma lifetime t_f to the initial plasma conditions and the phase transition temperature T_c

$$t_i T_i^3 = t_f T_c^3\,. \tag{6}$$

We thus arrive at a relation between plasma lifetime and energy in the plasma phase, valid for all nuclei

$$E/A^{2/3} = \tfrac{1}{15}\pi^3[3\alpha/(4b+3)]1.44lT_c^4$$
$$\times(t_f/t_i)^{4/3}[1-(t_i/t_f)^{(4b+3)/3b}]\,. \tag{7}$$

3. Transverse energy dependence of the survival probability

The observed J/ψ resonances are either directly formed resonances of $c\bar{c}$ pairs or come from the decay of χ resonances. Potential model calculations [9] and lattice simulations of the heavy quark potential [10] suggest that the screening length in the plasma phase is small enough to prohibit the formation of χ states already at T_c whereas a somewhat larger temperature may be needed to dissolve the J/ψ states in a Debye-screened Coulomb potential. In our previous analysis of the momentum dependence of J/ψ suppression we assumed that direct J/ψ formation is

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[2] For a recent review see ref. [6].

[3] The contribution of strange quarks to the energy density is still small close to T_c [7]. It will increase α at most by 20%.

[4] As long as the plasma lifetime is small compared to the size of the projectile nucleus the suppression rate is determined mainly by the plasma lifetime [4,8]. This is the case for the oxygen data in the transverse energy range experimentally accessible. Otherwise the finite size of the nucleus becomes the dominant parameter [3] and the suppression rate at large p_T will not scale anymore with $E/A^{2/3}$ but rather with $E/A^{1/3}$.

impossible for temperatures larger than $T_{J/\psi}=1.3T_c$ whereas χ resonance formation is already prohibited at $T_\chi=T_c$. We will use the same critical temperatures in our present analysis of the E_T dependence. Using $b=1/3$ and $\alpha=18.5$ eq. (7) becomes

$$E/A^{2/3}=7.6x(t_f/t_i)^{4/3}[1-(t_i/t_f)^{13/3}]\text{GeV}\,, \qquad (8)$$

with x defined as

$$x=(l/1\text{fm})(T_c/200\text{ MeV})^4\,. \qquad (9)$$

The arbitrary scale parameter x has been defined such that it is equal to the length of the rapidity interval Δy covered by the plasma in the central region for the canonical choice of the transition temperature ($T_c=200$ MeV) [6] and the longitudinal extension of the plasma region at time $t_i=1$ fm ($l=l_0\Delta y$, $l_0=1$ fm) [11]. In the present set-up of the NA38 experiment the transverse energy is distributed over two units of rapidity ($\Delta y=2$). The canonical choice thus would be $x=2$. However, l_0 and T_c are both not well determined [#5]. We thus have to consider x as a free parameter.

For a given plasma lifetime t_f we can now determine the survival probability as function of p_T as described in ref. [4] and, using eq. (8), we can relate this to the total plasma energy measured in units of lT_c^4. In fig. 1 we show the survival probability of J/ψ resonances with $p_T=0$ for $x=1.0$, 1.5 and 2.0. We note that there is a critical energy below which no suppression occurs. This critical value is related to the formation time τ_χ of χ resonances; for $t_f\leqslant\tau_\chi$ all normal χ-resonance formation is still possible [#6]. For the critical energy we thus obtain

$$E_c/A^{2/3}=7.6x(\tau_\chi/t_i)^{4/3}[1-(t_i/\tau_\chi)^{13/3}]\text{ GeV} \\ =18.1x\text{ GeV}\,. \qquad (10)$$

The experimental results for the transverse energy dependence of J/ψ suppression are presented as functions of E_T as measured by the NA38 calorimeter, which provides the total neutral plus 1/3 of the charged particle transverse energy [1]. We should rescale our energy scale by a factor $\frac{2}{3}\cdot\frac{5}{9}$ in order to be able to compare with the experimental data [#7]. To this aim we introduce a transverse energy scale, using eq. (8), as

$$E_T/A^{2/3}=2.8x(t_f/t_i)^{4/3}[1-(t_i/t_f)^{13/3}]\text{ GeV}\,. \quad (11)$$

In fig. 2 we show our result for the E_T dependence of the survival probability at $p_T=0$ for $x=1.2$. This is compared with experimental results for the p_T-integrated survival probability for oxygen and sulphur

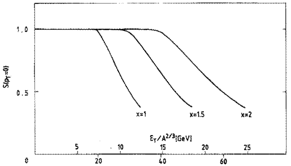

Fig. 1. Survival probability for J/ψ with $p_T=0$ versus $E/A^{2/3}$ for three different values of x. Also shown is the "experimental" transverse energy scale as defined in eq. (11).

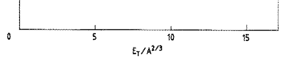

Fig. 2. Survival probability for J/ψ with $p_T=0$ versus $E_T/A^{2/3}$ for $x=1.2$. Also shown are the experimental data for the p_T-integrated survival probability for oxygen–uranium (●) and sulphur–uranium (■). Data are taken from ref. [12] and normalized to the lowest E_T-bin.

[#5] Note that T_c^4 enters the definition of x. Already a 20% uncertainty in the value of T_c amplifies to a factor 2 uncertainty in x.

[#6] We use $\tau_\chi=2.01$ fm [4]. For the suppression of direct J/ψ the plasma lifetime has to be even larger than τ_χ ($t_f\geqslant2.2$ fm) as we have assumed that the initial temperature has to be larger than $1.3T_c$ in order to affect these resonances.

[#7] The E_T values quoted in ref. [1] still have to be corrected for the calorimeter efficiency. This, however, has been taken into account in [12].

[12]. The experimental data have been normalized to the lowest E_T-bin. This should eliminate a great part of the suppression effects due to nuclear absorption which are expected to be only weakly dependent on E_T [13]. The plasma model reproduces quite well the slope of the E_T dependence for the large E_T events. For $x=1.2$ we find that there is no J/ψ suppression below $E_T/A^{2/3}=8.2$ GeV, i.e. below $E_T=52$ (82) GeV for oxygen (sulphur). For the canonical choice of the longitudinal extension of the plasma region, $l_0=1$ fm, we find that $x=1.2$ corresponds to a transition temperature $T_c=180$ MeV, an acceptable value given the tolerance in the other parameters entering the model (formation times $\tau_{J/\psi}$, τ_χ, ratio of critical temperatures $T_{J/\psi}/T_\chi$).

With the scale factor x we have fixed all parameters of the model. In particular we have fixed for any given value of $E_T/A^{2/3}$ the plasma lifetime and with that the p_T dependence of the suppression pattern. In fig. 3 we show the p_T dependence for five different values of $E_T/A^{2/3}$ using $x=1.2$. These transverse energy values have been chosen such that they fall into the different E_T-bins selected by the NA38 Collaboration for their oxygen data samples. The pattern shown in fig. 3 actually reproduces quite well the one observed by the NA38 collaboration.

In our present calculation we tacitly assumed that we are dealing with central collisions. However, at present the NA38 experiment does not trigger on these events but rather integrates over all impact parameters. Whereas the high E_T events mainly come from central events [8], those with low E_T are mainly peripheral events [14]. Our computation thus should be valid only at large E_T. In order to describe also the low E_T region with our ansatz it would be necessary to take into account the varying impact parameters per E_T-bin [13]. At present we overestimate the size of the plasma region for non-central events. A correct incorporation of the impact parameter dependence would thus lead to a flattening of our theoretical curve for small $E_T/A^{2/3}$.

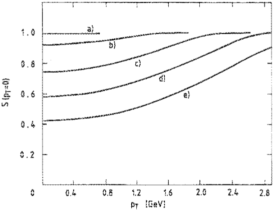

Fig. 3. Survival probability for J/ψ versus p_T for five different values of $E_T/A^{2/3}$. Shown are results for $E_T/A^{2/3}\leqslant 8.2$ GeV (a) and $E_T/A^{2/3}=9.4$ (b), 11.0 (c), 12.6 (d) and 14.2 (e) GeV.

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[8] In oxygen–uranium collisions events with $E/A^{2/3}\geqslant 9$–10 GeV are mainly due to central collisions [11].

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