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Production of η_c and dimuon enhancement in heavy ion collisions

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Abstract. Dilepton production in heavy ion collisions in the intermediate mass region (IMR) has consistently shown an excess over theoretical estimates. An attempt to understand this discrepancy between the observed dilepton pairs and the theoretical estimate is made here through the production of the η_c meson and estimates obtained by NRQCD calculations. We find that η_c production offers a satisfactory quantitative picture for explaining the discrepancy.

Dilepton production plays a very important role in the study and understanding of heavy ion collisions. This is mainly because dileptons do not interact with the surrounding hadronic medium after being produced in a nucleus–nucleus collision. The typical dilepton invariant mass spectrum appears as a wide continuum interrupted by various resonance ($\rho, \omega, \phi, J/\psi, \psi'$ etc) decay peaks.

The intermediate mass region (between the ϕ and J/ψ peaks, 1.5 GeV to 2.5 GeV) is particularly interesting because it is believed to contain dileptons created in the thermalized QGP produced in nucleus–nucleus collisions (thermal dimuons). However, it is precisely in this region that many different experiments in A–A collisions [1] have shown an excess of dimuon production. In all these data sets, the dilepton sources are either Drell–Yan pairs or decays of J/ψ or $D\bar{D}$. While the bulk of the data agree with such a production picture, there is a significant discrepancy between the observed dilepton pairs and the theoretical estimate based on the above sources in this intermediate mass region (IMR) with a $\mu^+\mu^-$ invariant mass in the range 1.5–2.5 GeV.

Many explanations have been offered in the literature for this excess, viz. a decrease in the ρ meson mass due to thermal effects in e^+e^- data [2], D-rescattering [3], enhanced $D\bar{D}$ production, in-flight $\pi^+\pi^-$ decaying to e^+e^- [4], fireball hydrodynamics [5] and so on.

In the entire kinematic range (up to 5 GeV) other than the IMR region mentioned above, where theory and experiment agree, the overall picture involving charm quarks [6] is that when there is sufficient energy exchange in a collision, protons have a non-negligible charm content ($c\bar{c}$ pairs), and substantial high energy gluons which in turn can decay to $c\bar{c}$ pairs. These $c\bar{c}$ pairs occasionally form bound states such as J/ψ by emitting a soft gluon to maintain color balance, or they can further polarize $u\bar{u}$ or $d\bar{d}$ from the surrounding medium to form $D\bar{D}$ pairs. Although the present theoretical understanding cannot predict absolute numbers for these processes, it is possible to check the consistency of this picture with various p–A and A–A data. By and large the data agree with various quantitative checks.

Other charm meson bound states can also be produced, such as η_c, ψ' and χ''s, though their relative abundance is constrained by size, the larger ones being less likely to be formed.

However, the η_c meson, which is a $1S$ orbital state, is expected to have the same size and has almost the same mass as the J/ψ. They only differ in their spin and hence in any collision where $c\bar{c}$ quarks are produced these can form η_c with about 1/3 probability compared to J/ψ. However, the possibility of the production of η_c mesons in nucleus–nucleus collisions was only considered very recently in [7] wherein this was presented as a possible explanation of the discrepancy between theory and experiment in nucleus–nucleus collisions – in particular in S–U and Pb–Pb collisions. However, this paper presented a somewhat qualitative field theoretical picture of the production of η_c mesons and used it to explain the discrepancy in the IMR region. Furthermore, this simple field theoretical picture was unable to account for the discrepancy seen in the IMR region for Pb–Pb collisions in the central region.

In what follows we attempt to understand the discrepancy in the IMR region once again through the production of the η_c, but by using a more quantitative description for the production of the η_c meson. Our attempt may, therefore, be seen as a way to build upon the ideas of [7] using the technology of NRQCD to present a more quantitative picture.

In NRQCD [8, 9], the quarkonium production cross section factorizes into a perturbatively calculable short-distance ($\leq 1/M_Q$, M_Q is the mass of the heavy quark) effect and a long-distance part which is given by non-

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perturbative matrix elements. The cross section for the production of a quarkonium state H can be written as

$$\sigma(H) = \sum_n \frac{F_n}{M_Q^{d_n-4}} \left\langle 0 \left| \mathcal{O}_n^H \right| 0 \right\rangle . \tag{1}$$

The coefficients F_n correspond to the production of $Q\overline{Q}$ in the angular momentum and color state (singlet or octet) denoted by n and is calculated using perturbative QCD. The non-perturbative part, $\langle \mathcal{O}_n^H \rangle$, of mass dimension d_n in NRQCD, has a well-defined operator definition and is universal. These matrix elements can be extracted from any one process and can then be used to predict other processes where the same matrix elements appear. Though the summation involves an infinite number of terms, the relative magnitude of the various terms is predicted by NRQCD and these matrix elements scale as powers of the relative velocity v. However, this does not necessarily imply that effects from higher orders in v will always be small in physical processes, because any observable, like the decay width or the cross section, is given by a double expansion in the strong coupling constant $\alpha_s(M_Q)$ and the relative velocity v.

The non-perturbative matrix elements in NRQCD are not calculable and have to be obtained by fitting to the available data. The matrix elements of the color-singlet operators can be obtained from the quarkonium decay widths, and the color-octet matrix elements have been obtained by fitting the NRQCD predictions to the CDF data [10, 11]. The remarkable thing is that the non-perturbative parameters appearing in the η_c production cross section can be determined from the matrix elements determined from J/ψ production at the Tevatron: this happens because of the heavy-quark symmetry of the NRQCD Lagrangian. This has been exploited earlier in the context of h_c and η_c production at Tevatron [12, 13].

A Fock space expansion of the physical η_c, which is a ${}^1S_0$ ($J^{PC} = 0^{-+}$) state, yields

$$\begin{aligned} |\eta_c\rangle = \mathcal{O}(1) \left| Q\overline{Q} \left[{}^1S_0^{[1]} \right] \right\rangle + \mathcal{O}(v^2) \left| Q\overline{Q} \left[{}^1P_1^{[8]} \right] g \right\rangle \\ + \mathcal{O}(v^4) \left| Q\overline{Q} \left[{}^3S_1^{[8]} \right] g \right\rangle + \ldots \end{aligned} \tag{2}$$

The color-singlet ${}^1S_0$ state contributes at $\mathcal{O}(1)$ but the color-octet ${}^1P_1$ and ${}^3S_1$ channels effectively contribute at the same order, because the P-state production is itself down by a factor of $\mathcal{O}(v^2)$. The color-octet states become a physical η_c by the ${}^1P_1^{[8]}$ state emitting a gluon in an E1 transition, and by the ${}^3S_1^{[8]}$ state emitting a gluon in an M1 transition. The contributing subprocess cross sections are

$$\begin{aligned} q\ \bar{q} &\to Q\overline{Q} \left[{}^{2S+1}L_J \right] , \\ g\ g &\to Q\overline{Q} \left[{}^{2S+1}L_J \right] , \end{aligned}$$

where the $Q\overline{Q}$ is in the ${}^1S_0^{[1]}$, ${}^1S_0^{[8]}$ and ${}^3S_1^{[8]}$ states. The ${}^1P_1^{[8]}$ state does not contribute at this leading order in α_s because of Yang's theorem.

We have computed the contributions to the cross section for η_c production from the ${}^1S_0^{[1]}$, ${}^1S_0^{[8]}$ and ${}^3S_1^{[8]}$ states.

Heavy-quark spin symmetry is made use of in obtaining the $\langle \mathcal{O}_n^{\eta_c} \rangle$ from the experimentally available $\langle \mathcal{O}_n^{J/\psi} \rangle$. Using this symmetry we get the following relations among the $\langle \mathcal{O}_n^H \rangle$:

$$\begin{aligned} \langle 0 | \, \mathcal{O}_1^{\eta_c} \left[{}^1S_0 \right] |0\rangle &= \langle 0 | \, \mathcal{O}_1^{J/\psi} \left[{}^3S_1 \right] |0\rangle \ (1 + O(v^2)), \\ \langle 0 | \, \mathcal{O}_8^{\eta_c} \left[{}^3S_1 \right] |0\rangle &= \langle 0 | \, \mathcal{O}_8^{J/\psi} \left[{}^1S_0 \right] |0\rangle \ (1 + O(v^2)). \end{aligned} \tag{3}$$

For the singlet matrix elements we have $\langle 0 | \, \mathcal{O}_1^{J/\psi} [{}^3S_1] \, |0\rangle = 1.2\,\mathrm{GeV}^3$. The CDF J/ψ data only constrain a combination of octet matrix elements given by $A_1 + A_2 \equiv \frac{\langle 0|\mathcal{O}_8^{J/\psi}[{}^3P_0]|0\rangle}{M_c^2} + \frac{\langle 0|\mathcal{O}_8^{J/\psi}[{}^1S_0]|0\rangle}{3} = (2.2 \pm 0.5) \times 10^{-2}\,\mathrm{GeV}^3$ [11]. The CDF J/ψ data do not allow for a separate determination of the values of A_1 and A_2, because the shapes of these two contributions to the J/ψ p_T distribution are almost identical. For our numerical predictions we assume that the non-perturbative matrix element of interest to us ($\langle 0 | \, \mathcal{O}_8^{\eta_c} \left[{}^3S_1 \right] |0\rangle$) lies in the range determined by this sum.

Using these values for the non-perturbative matrix elements we can compute the cross section for η_c production. This has to be convoluted with the branching ratio for $\eta_c \to \gamma\gamma^* \to \gamma\mu^+\mu^-$. This quantity can be estimated only approximately, and to do this we have used the value of the $\eta_c \to \gamma\gamma$ branching ratio modulated by a $(1 - M_{\gamma^*}^2 / M_{\eta_c}^2)$ in the denominator. With these inputs, we have computed the $\mu^+\mu^-\gamma$ yield in Pb–Pb and S–U collisions. The nuclear parton distributions have been taken from the parameterisation of EKS [14]. The cuts and acceptances for the muons and photons have been taken from the experimental papers (see, for example, [15]).

Our results for $\mathrm{d}N/\mathrm{d}M$ are compared to the experimental curves for central collisions in Fig. 1. In keeping with the convention used by the experimentalists, we have indicated the η_c production contribution separately (long-dashed line) as well as the summed contribution of all the individual contributions in the IMR region (dashed-dotted line). It appears from both graphs that the η_c contribution saturates the discrepancy seen in the IMR region. However, this agreement is based on certain assumptions. First of all, the branching ratio of $\eta_c \to \gamma\gamma^* \to \gamma\mu^+\mu^-$ used, as already stated in the previous paragraph, is only approximately estimated through very general field theoretical arguments. Secondly, as with all leading order calculations, the result is dependent (though mildly) on the scale of α_s used in the calculation. Finally, since, as we have explained, the CDF J/ψ data only constrain the combination of octet matrix elements given by A_1 and A_2, we have, for convenience, assumed that the non-perturbative matrix element $\langle 0|\mathcal{O}_8^{J/\psi}[{}^1S_0]|0\rangle$ (i.e. A_2) saturates the sum and hence gives the value of the $\langle 0|\mathcal{O}_8^{\eta_c}[{}^3S_1]|0\rangle$ matrix element. However, this turns out to be not such a serious approximation, and we find that varying this value between A_1 and A_2 produces very little change in the overall shape and magnitude of the η_c contribution.

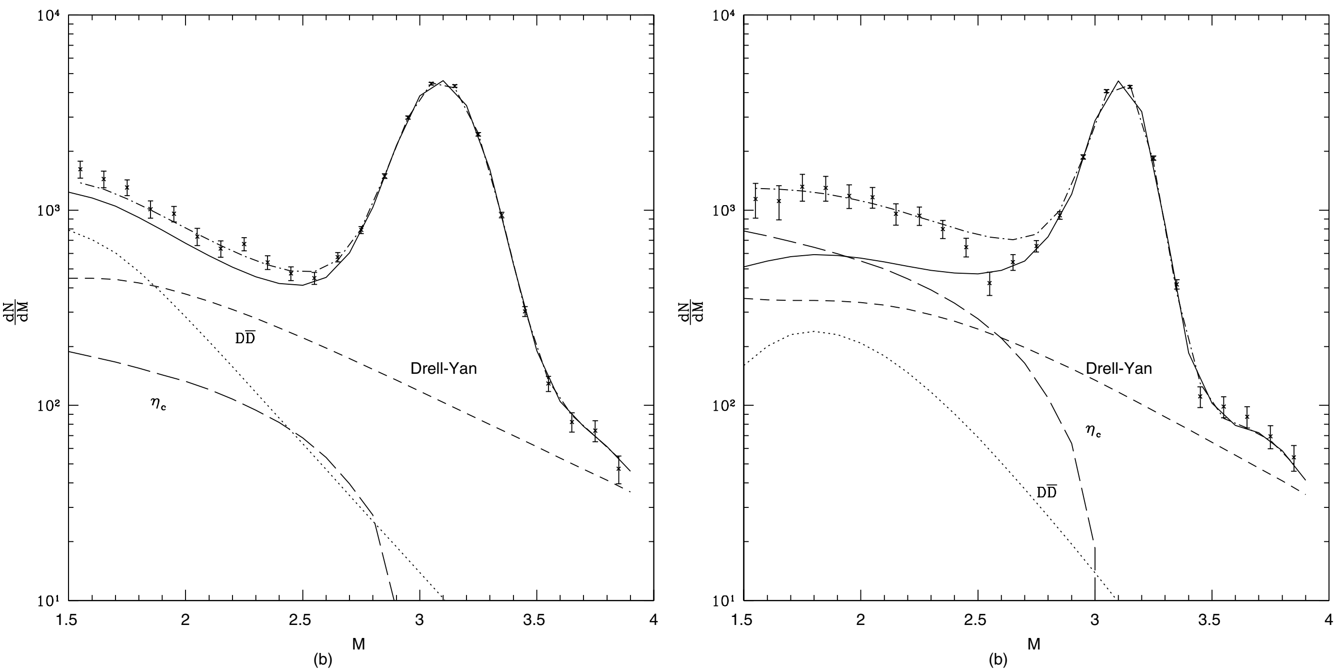

Fig. 1. The graphs show the S–U (a) and Pb–Pb (b) results for central collisions. Note that in the IMR region the η_c contribution appears to saturate the deficit

The general arguments for the η_c contribution used in [7] were unable to explain the discrepancy in central collisions for in Pb–Pb, and it was speculated there that the difference between theory and experiment even after including η_c production could perhaps be accounted for by glueball production. It seems, however, in the present, more quantitative, analysis, that this possibility is ruled out since the dileptons from η_c seem to cover the deficit. In any case, even if such a scenario of glueball production is envisaged, the contribution would clearly be very small.

Finally, we have made no comments about peripheral collisions where also this discrepancy in the IMR region is seen. This is because a quantitative study of peripheral collision would involve a detailed understanding of the geometry of the collision, and its consequent uncertainties. In view of these, we feel that no firm statements can really be made within the context of this approach to peripheral collisions in heavy ion collisions.

Our calculation based on NRQCD will probably get significantly corrected via non-perturbative corrections at the low values of $p_{\rm T}$ that we are considering. Moreover, the cross section is also vulnerable to the usual perturbative uncertainties characteristic of leading order QCD calculations. Another source of uncertainty is the η_c branching ratio to γe^+e^- which is not available from the data. Given these limitations, the numbers presented in the paper must be treated as ball-park figures which, nevertheless, constitute the first quantitative estimate of the η_c contribution to the dilepton yield and will hopefully spur both experimental interest and more sophisticated theoretical efforts.

In conclusion, in this paper we have used the technology of NRQCD to estimate the contribution of the η_c meson to dilepton production in the IMR region of heavy ion collisions and found that it is able to satisfactorily explain the deficit, within the limits of the assumptions already mentioned earlier.

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