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 $\eta_c$ meson and estimates obtained by NRQCD calculations. We find that $\eta_c$ 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 $\phi$ and $J/\psi$ peaks, 1.5 GeV to 2.5 GeV) is particularly interesting because it is believed to contain dileptons created in the thermalised QGP produced in nucleus-nucleus collisions (thermal dimuons). However it is precisely in this region that many different experiments in A-A collisions \cite{1} have shown an excess of dimuon production. In all these data sets, the dilepton sources are either Drell-Yan pairs or decays of $J/\psi$ or $D\bar{D}$. While the bulk of the data agrees 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 $\mu^+\mu^-$ invariant mass in the range 1.5 – 2.5 GeV.

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

In the entire kinematic range (upto 5 GeV) other than the IMR region mentioned above, where theory and experiment agree, the overall picture involving charm quarks \cite{6} is that when there is sufficient energy exchange in a collision, protons have 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 $J/\psi$ by emitting a soft gluon to maintain color balance, or 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 agrees with various quantitative checks.

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

However the $\eta_c$ meson, which is a 1S orbital state is expected to have the same size and has almost the same mass as the $J/\psi$. They only differ in their spin and hence in any collision where $c\bar{c}$ quarks are produced these can form $\eta_c$ with about 1/3 probability as $J/\psi$. However the possibility of the production of $\eta_c$ mesons in nucleus-nucleus collisions was only considered very recently in \cite{7} wherein this was presented as a possible explanation of the discrepancy between theory and experiment in nucleus-nucleus collision – in particular in S-U and Pb-Pb collisions. However this paper presented a somewhat qualitative field theoretical picture.
of the production of $\eta_c$ 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 $\eta_c$ but by using a more quantitative description for the production of the $\eta_c$ 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 factorises into a perturbatively calculable short-distance ($\lesssim 1/M_Q$, $M_Q$ is the mass of the heavy quark) effect and a long-distance part which is given by non-perturbative matrix elements. The cross section for production of a quarkonium state $H$ can be written as

$$\sigma(H) = \sum_n \frac{F_n}{M_Q^{d_n-1}} \langle 0 | O^H_n | 0 \rangle.$$  

(1)

The coefficients $F_n$ correspond to the production of $Q\bar{Q}$ in the angular momentum and colour state (singlet or octet) denoted by $n$ and is calculated using perturbative QCD. The non-perturbative part, $\langle O^H_n \rangle$ 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 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 available data. The matrix elements of the colour-singlet operators can be obtained from the quarkonium decay widths and the colour-octet matrix elements have been obtained by fitting NRQCD predictions to the CDF data [10, 11]. The remarkable thing is that the non-perturbative parameters appearing in the $\eta_c$ production cross section can be determined from the matrix elements determined from $J/\psi$ 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 $\eta_c$ production at Tevatron [12, 13].

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

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

The colour-singlet $1S_0$ state contributes at $\mathcal{O}(1)$ but the colour-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 colour-octet states become a physical $\eta_c$ 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

$$q \bar{q} \to Q\overline{Q}[2^{S+1}L_J],$$
$$g g \to Q\overline{Q}[2^{S+1}L_J],$$

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 $\alpha_s$ because of Yang’s theorem.

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

Heavy quark spin-symmetry is made use of in obtaining $\langle \mathcal{O}_{\eta_c}^{[8]} \rangle$’s from the experimentally available $\langle \mathcal{O}_{\eta_c}^{[1]} \rangle$’s. Using this symmetry we get the following relations among $\langle \mathcal{O}_{\eta_c}^{[8]} \rangle$’s:

$$\langle 0 | O_8^{\eta_c} [1S_0] | 0 \rangle = \langle 0 | O_1^{J/\psi} [3S_1] | 0 \rangle (1 + O(v^2)),$$
$$\langle 0 | O_8^{\eta_c} [3S_1] | 0 \rangle = \langle 0 | O_1^{J/\psi} [1S_0] | 0 \rangle (1 + O(v^2)). \tag{3}$$

For the singlet matrix elements we have $\langle 0 | O_1^{J/\psi} [3S_1] | 0 \rangle = 1.2$ GeV$^3$. The CDF $J/\psi$ data only constrains a combination of octet matrix elements given by $A_1 + A_2 \equiv \langle 0 | O_8^{\eta_c} [3P_0] | 0 \rangle + \langle 0 | O_8^{\eta_c} [1S_0] | 0 \rangle = (2.2 \pm 0.5) \times 10^{-2}$ GeV$^3$ \cite{11}. The CDF $J/\psi$ 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 $J/\psi p_T$ distribution are almost identical. For our numerical predictions we assume that the non-perturbative matrix element of interest to us ($\langle 0 | O_8^{\eta_c} [3S_1] | 0 \rangle$) lies in the range determined by this sum.

Using these values for the non-perturbative matrix elements we can compute the cross-section for $\eta_c$ production. This has to be convoluted with the branching ratio for the $\eta_c \to \gamma\gamma^* \to \gamma\mu^+\mu^-$. 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 \cite{13}. The cuts and acceptances for the muons and photons have been taken from the experimental papers (see, for example, \cite{15}).

3
Figure 1: The graphs show the S-U (a) and Pb-Pb (b) results for central collisions. Note that in the IMR region the $\eta_c$ contribution appears to saturate the deficit.

Our results for $dN/dM$ 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 $\eta_c$ 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 the graphs that the $\eta_c$ 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 $\alpha_s$ used in the calculation. Finally since, as we have explained, the CDF $J/\psi$ data only constrains the combination of octet matrix elements given by $A_1$ and $A_2$ we have, for convenience, assumed that the non-perturbative matrix element $\langle 0|\mathcal{O}^{J/\psi}_{8}[1S_0]|0 \rangle$ (i.e $A_2$) saturates the sum and hence gives the value of the $\langle 0|\mathcal{O}^{J/\psi}_{8}[3S_1]|0 \rangle$ 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 $\eta_c$ contribution.

The general arguments for the $\eta_c$ contribution used in [7] were unable to explain the discrepancy in central collision in Pb-Pb collision and it was speculated there that the difference between theory and experiment even after including $\eta_c$ produc-
tion, 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 $\eta_c$ 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.

In conclusion, in this paper, we have used the technology of NRQCD to esti-
mate the contribution of the $\eta_c$ 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 used in this study.
References

[1] I. Ravinovich for CERES Coll., Nucl. Phys. A638, 159C (1998); G. Agakichiev for CERES Coll., Nucl. Phys. B422, 405 (1998); M. Masera for HELIOS3 Coll., Nucl. Phys. A590, 93C (1995); A. De Falco for NA38 Coll., Nucl. Phys. A638, 487C (1998); E. Scomparin for NA50 Coll., J. Phys. G25, 235 (1999); A. Drees, Nucl. Phys. A610, 536C (1996).

[2] G. Q. Li, C. M. Ko and G. E. Brown, Phys. Rev. Lett. 75, 4007 (1995); Nucl. Phys. A606, 568 (1996).

[3] Z. Lin and X. N. Wang, Phys. Lett. B444, 245 (1998).

[4] I. Tserruya, LANL Archives nucl-ex/9912003 and references therein.

[5] R. Rapp and E. Shuryak, Phys. Lett. B473, 13 (2000).

[6] D. Kharzeev, Nucl. Physics A638, 279C (1998) and references therein.

[7] R. Anishetty and R. Basu, Phys. Lett. B495, 295 (2000).

[8] G.T. Bodwin, E. Braaten and G.P. Lepage, Phys. Rev. D51, 1125 (1995); erratum ibid. 55 5853 (1997).

[9] M. Krämer, Prog. Part. Nucl. Phys. 47 141 (2001).

[10] P. Cho and A.K. Leibovich, Phys. Rev. D53, 150 (1996).

[11] P. Cho and A.K. Leibovich, Phys. Rev. D53, 6203 (1996).

[12] K. Sridhar, Phys. Rev. Lett., 77, 4880 (1996).

[13] P. Mathews, P. Poulose and K. Sridhar, Phys. Lett. B438 336 (1998).

[14] K.J. Eskola, V.J. Kolhinen and P.V. Ruuskanen, Nucl. Phys. B 535 (1998) 351.

[15] M. C. Abreu et al., CERN-EP-2000-012, Eur. Phys. J. C14, 443 (2000).