Nuclear Broadening of Transverse Momentum in Drell-Yan Reactions

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Data for Drell-Yan (DY) processes on nuclei are currently available from fixed target experiments up to the highest energy of \(\sqrt{s} = 40\) GeV. The bulk of the data cover the range of short coherence length, where the amplitudes of the DY reaction on different nucleons do not interfere. In this regime, DY processes provide direct information about broadening of the transverse momentum of the projectile parton experiencing initial-state multiple interactions. We revise a previous analysis of data from the E772 experiment and perform a new analysis of broadening including data from the E866 experiment at Fermilab. We conclude that the observed broadening is about twice as large as the one found previously. This helps to settle controversies that arose from a comparison of the original determination of broadening with data from other experiments and reactions.

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INTRODUCTION

Understanding how partons propagate in a nucleus reveals precious information about the early stages of hadronization that is difficult to obtain by other means. This understanding also plays a role in using partons as a sensitive probe of matter produced in relativistic heavy ion collisions. One of the important observables is broadening of the transverse momentum of produced particles due to multiple interactions of a parton in the medium. This effect can be studied for the nuclear medium in Drell-Yan reactions \[1,2,3\] and deep-inelastic scattering \[4,5\] on nuclear targets. The former is an especially clean tool, since the produced dilepton, which has no final state interactions, carries undisturbed information on the transverse momentum of the parton initiating the reaction.

A parton propagating through a nuclear medium experiences Brownian motion in the transverse momentum frame \[6\], with the increase of the mean transverse momentum squared being proportional to the length of the path times the medium density \(\rho(z)\),

\[
\Delta \langle p_T^2(L) \rangle = 2C(s) \int_0^L dz \rho(z),
\]

where \(z\) is the longitudinal coordinate. The energy dependent dimensionless factor \(C\) characterizes multiple interactions of the parton with bound nucleons. In the light-cone dipole description of the broadening process \[7,8,9\] \(C\) is related to the small separation behavior of the dipole-nucleon cross section,

\[
C(s) = \left. \frac{d\sigma_{qq}(r_T,s)}{dr_T^2} \right|_{r_T=0}.
\]

The cross section \(\sigma_{qq}(r_T,s)\) introduced in \[10\] varies with transverse \(qq\) separation \(r_T\) and energy. This quantity is difficult to calculate, since it involves nonperturbative effects, but it can be fitted to data for the photoabsorption cross section and the proton structure function \(F_2(Q^2,x)\), which probe a wide variety of separations and energies \[11,12\]. Since each of the scatterings is accompanied by a loss of energy, there is a close connection between the multiple scattering that leads to momentum broadening and initial state gluonic energy loss \[13,14\]. It was found, however, that this induced energy loss effects the DY process much less than the vacuum gluon radiation initiated by the interaction on the nuclear surface \[13\].

In the parton model, a single interaction of a quark propagating through the medium can be viewed as a quark-gluon correlation in a nucleus in its infinite momentum frame. The corresponding twist-4 term \[10\] cannot be calculated, but is rather fitted to data. This phenomenon has not been successful so far, with the parameter extracted from data varying by a factor of 20, depending on assumptions \[17\] and also demonstrating a strong process dependence \[18\].

The first observation of nuclear enhancement of large transverse momenta was made in hadron-nucleus collisions \[13\]. A sensitive tool to measure the transverse momentum broadening of quarks propagating through nuclei is heavy dilepton production in the Drell-Yan reaction. The leptons have no interaction in the nuclear medium and thus carry undisturbed information about the transverse momentum of the quark. Nuclear broadening of the mean transverse momentum squared of the...
dilepton is defined as
\[ \Delta(p_T^2) = \langle p_T^2 \rangle_A - \langle p_T^2 \rangle_N . \] (3)
Perturbative QCD predicts \( d\sigma_{DY}/dp_T^2 \propto 1/p_T^4 \) (for transversely polarized virtual photons) at large \( p_T \). Therefore, the mean value \( \langle p_T^2 \rangle \) is divergent, and the broadening of Eq. (3) might be ill defined (see discussion in \[20\]). Of course, the experiment imposes a cut off. Thus, the experimentally observed broadening is finite, but it may depend on the experimental acceptance. This divergence, however, sometimes cancels in the broadening for several reasons. First of all, perturbative QCD predicts no nuclear effects at large \( p_T \), i.e., \( d\sigma_{DY}(pA)/dp_T^2 \rightarrow A d\sigma_{DY}(pp)/dp_T^2 \) (see, e.g., \[21\]). Secondly, nuclear shadowing is expected to vanish at large \( p_T \), especially at an energy as low as 800 GeV, the energy of our current analysis. In this case, the divergent high-\( p_T \) part of the integrations cancel in (3), and one can measure the nuclear broadening effect within a restricted interval of \( p_T \) where the nuclear effects are visible.

An analysis of broadening in the Drell-Yan reaction was presented in \[3\] using data from the E772 experiment at Fermilab. The results of the analysis for \( \Delta(p_T^2)_{DY} \) are depicted in Fig. 1 as function of \( A \).

![FIG. 1: Broadening of transverse momentum in production of heavy dileptons (squares), J/Ψ (circles) and Υ (triangles) in \( pA \) collisions at 800 GeV. The data are from the E772, E789, and E866 experiments at Fermilab \[2, 22, 23, 34\]. The curves are fit results with the parametrization \( \Delta(p_T^2) = D [(A/2)^{1/3} - 1] \).](image)

Following Eq. (1), one can parametrize the \( A \) dependence as
\[ \Delta(p_T^2) = D [(A/2)^{1/3} - 1] \] (4)
A fit to the data points for the Drell-Yan reaction in Fig. 1 results in \( D = 0.027 \text{GeV}^2 \), which corresponds to \( C = 2 \) in (1).

The same figure presents nuclear broadening in \( J/Ψ \) and \( Υ \) production at the same energy. Comparing with this data and results available at lower energies, one can make the following observations: (i) the magnitude of broadening in reactions of heavy dilepton and quarkonium production differ by a factor of 5, while one would expect the Casimir factor 9/4. Indeed, heavy quarkonia should be produced via gluons, which interact in the medium more strongly than quarks. (ii) The errors for Drell-Yan data are small compared to \( J/Ψ \) production, whereas the latter is measured with much higher statistics. (iii) Comparison with broadening measured at lower energies presented in Fig. 2 suggests that broadening in heavy quarkonium production should rise with energy, whereas it is seen to be rather flat, or even perhaps drop, at high energies in the Drell-Yan reaction.

![FIG. 2: Energy dependence of the parameter \( D \) in (4) fitted to data for dilepton and heavy quarkonium production in \( pA \) and \( πA \) collisions.](image)

The latter observation is difficult to explain, since the cross section controlling the multiple interactions in the medium rises with energy. (iv) The dipole approach, Eq. (2), based on a phenomenological color dipole cross section fitted to data, predicts the factor \( C \) \[8\] to be about twice as large as one suggested by the data presented in Fig. 1.

In this paper we revise the result of the previous analysis \[3\] of the E772 data for Drell-Yan nuclear ratios, combining them with the brand-new data \[25\] from the E866 experiment for Drell-Yan cross section in \( pD \) collisions. The improvement in the analysis arises from the fact that the E866 experiment provides more accurate \( pD \) cross section data at small \( p_T \) than that available from E772. Additionally, it was assumed with no good justification that the same simple parametrization of the \( p_T \) dependence is valid for both \( pD \) and \( pA \) collisions, which dictated a specific shape for the ratios and led to a sub-
stantial underestimate of the errors for \( p_T \)-broadening. Our new analysis is less model-biased, but it results in somewhat larger errors. We also perform here an analysis of new data from the E866 experiment, which leads to a broadening about twice as large as in [3], although with larger experimental errors. This result resolves the problems stated earlier.

**LIGHT-CONE DIPOLE DESCRIPTION OF NUCLEAR BROADENING**

When a projectile scatters from a nucleus, one can treat the target nucleons as independent with some degree of accuracy if viewed in the rest frame of the nucleus, since the bound nucleons are sufficiently well separated in that frame. However, when boosted to the infinite momentum frame the same reaction looks different since the nucleons may start communicating via their parton clouds at small Bjorken \( x \). Indeed the Lorentz contraction is proportional to \( 1/x \), therefore it is much stronger at large, than at small \( x \). As a result, the parton clouds are still separated for \( x > x_c \), but overlap and interact at smaller \( x, x < x_c \). The border line \( x_c \) can be estimated as,

\[
x_c \sim \frac{1}{m_N R_A},
\]

where \( R_A \) is the nuclear radius. Note that this is only an order of magnitude estimate. There is a numerical factor in [51] that varies considerably for different species of partons and the way the parton distribution is probed [27].

The interaction between partons at \( x < x_c \) leads to a modification of their transverse momentum distribution: suppression of small momenta, enhancement at intermediate, and no effect at large transverse momenta. This phenomenon is known as the saturation effect [28], or color glass condensate [29]. One should keep in mind, however, that onset of this effect happens only at \( x < x_c \).

The partonic interpretation of hard reactions at small Bjorken \( x \) is known to be Lorentz non-invariant. Deep-inelastic scattering (DIS) looks like absorption of the virtual photon by a quark (antiquark) that belongs to the target in the infinite momentum frame of the latter. The reaction looks, however, quite different in the target rest frame, as the interaction with a \( \bar{q}q \) fluctuation of the high-energy virtual photon. Of course, all observables, including the structure function \( F_2(x, Q^2) \), which is proportional to the total photon-target cross section, are Lorentz invariant.

The Drell-Yan reaction is usually interpreted as annihilation of a quark with an antiquark belonging to the beam and target respectively (or vice versa), \( \bar{q}q \rightarrow \gamma^* \) [30]. While this is correct in the rest frame of the dilepton, this interpretation of a heavy dilepton production is not applicable in target rest frame. Indeed, in order to satisfy the kinematics annihilating with the projectile, the target parton must move with momentum \( p_L \sim p_T/x_2 \), where \( x_2 \) in standard notation is the Bjorken \( x \) of the target parton. At small \( x_2, x_2 \ll 1 \), this would mean presence of highly energetic partons in the stationary target. The proper space-time interpretation in this case is analogous to the Weizsäcker-Williams mechanism of electromagnetic bremsstrahlung. Namely, a beam quark (antiquark) develops a fluctuation, \( q \rightarrow q \gamma^* \) (with \( \gamma^* \rightarrow \gamma \)), which interacts with the target freeing the dilepton [31, 32].

What looked like overlap of parton clouds at small \( x_2 \) in the infinite momentum frame, now takes the form of interference between the amplitudes of a parton interacting with different nucleons. On the same footing it can also be interpreted as the lifetime of the \( |q \gamma^* \rangle \) fluctuation, given by the uncertainty principle

\[
l_c = \frac{1}{q_L} = \frac{2E_q}{M_{\gamma^*}^2 - m_q^2},
\]

where \( q_L \) is the longitudinal momentum transfer in the \( q \rightarrow q \gamma^* \) fluctuation. The effective mass squared of the fluctuation is \( M_{\gamma^*}^2 = M_{\gamma^*}^2/\alpha + m_q^2/(1-\alpha) + k_T^2/\alpha(1-\alpha) \), where \( \alpha \) is the fraction of the light-cone momentum of the parent quark taken by the dilepton (\( \gamma^* \)), and \( k_T \) is the \( \gamma^* - q \) relative transverse momentum. Hereafter we do not discriminate between coherence time and length, assuming that the quark has the speed of light. Evidently, the same length \( \sim 1/q_L \) defines the maximal longitudinal distance between two scattering centers that can interfere. This is why the distance, Eq. (6), is also called coherence length.

In the case of a nuclear target, the coherence length is of great importance. If this time interval is small compared to the mean spacing of nucleons in the nucleus, dileptons are radiated in the Bethe-Heitler regime, i.e., with no interference between the amplitudes of radiation from different nucleons. At the another extreme, \( l_c \gg R_A \), the Landau-Pomeranchuk phenomenon governs the radiation. Namely, interferences suppress radiation with small transverse momenta, enhance it at intermediate values of few GeV, and make no changes at higher momentum transfer. In both cases, the linear dependence on the path length, or nuclear thickness, Eq. (1), holds, although with different factors \( C \) [20]. The borderline between the two regimes, \( l_c = R_A \), corresponds to the relation for \( x_c \), Eq. (5).

**Short coherence length**

The mechanism of nuclear broadening depends on how long the coherence length is compared to the nuclear size. In the regime of short coherence length, \( l_c \ll R_A \), the dilepton fluctuation appears for a very short time deep
inside the nucleus and is immediately released on mass shell by the interaction, which is assumed to be the same as on a free nucleon. Therefore, this interaction does not lead to any change in the transverse momentum dependence, i.e., it provides no broadening. However, the initial state interactions of the projectile quark affect the transverse momentum distribution of the quark, which acquires the shape [8]

\[ \frac{dN_{l\bar{l}}}{dp_T^2} = \int d^2r_1 d^2r_2 e^{i\vec{r}_1 \cdot \vec{r}_2} \Omega_{in}^q(\vec{r}_1,\vec{r}_2) \times e^{-\frac{1}{2}r_1^2+r_2^2} \Omega_0(\vec{r}_1,\vec{r}_2) \times T_A(b). \]  

(7)

Here \( \Omega_{in}^q(\vec{r}_1,\vec{r}_2) \) is the density matrix describing the impact parameter distribution of the quark in the incident hadron. We assume a Gaussian shape,

\[ \Omega_{in}^q(\vec{r}_1,\vec{r}_2) = \frac{\langle p_0^2 \rangle}{\pi} e^{-\frac{1}{2}(r_1^2+r_2^2)} \frac{\Omega_0}{(\vec{r}_1^2+\vec{r}_2^2)}, \]

(8)

where \( \langle p_0^2 \rangle \) is the mean value of the primordial transverse momentum squared of the quark.

The transverse momentum of the radiated dilepton reflects the broadening acquired by the quark in the nuclear medium. However, the effect is weaker than for the quark, since the transverse momenta of the dilepton and quark are connected by the relation,

\[ q_T^l = \alpha p_T^q, \]

(9)

where \( \alpha < 1 \) is the fractional momentum of the \( l\bar{l} \). This relation clearly demonstrates that the effect of broadening cannot be translated into a modification of the quark distribution function (as is frequently done), in the regime of short coherence length. Thus, it confirms the breakdown of \( k_T \)-factorization imposed by the initial state interaction, as suggested in [33].

No-shadowing sum rule

To the extent that the coherence length for the Drell-Yan reaction is short compared to the internucleon spacing in the nucleus, no shadowing is possible. This means that in spite of the nuclear modification of the transverse momentum dependence of Drell-Yan cross section, the number of quarks is conserved. This fact can be represented in the form of a sum rule,

\[ \int d^2p_T \frac{d\sigma_{DY}(pA)}{d^2p_T} = A \int d^2p_T \frac{d\sigma_{DY}(pN)}{d^2p_T}. \]

(10)

The nuclear ratios measured in the E772 and E866 experiments at 800 GeV depicted in Fig. 3 have a tendency to satisfy this constraint. Indeed, the ratios are below one at small \( p_T \), but rise above one at larger momentum transfers. The weak nuclear suppression observed at small \( p_T \), where the absolute cross section is maximal, seems to be compensated by a stronger enhancement at larger \( p_T \) in accordance with the sum rule Eq. (10). Data in the small \( p_T < 1 - 2 \) GeV domain, where the main part of the cross section comes from, are most important for broadening. In this region the E866 data in Fig. 3 demonstrate stronger nuclear suppression than the E772 data. The observed deviation from unity at small \( p_T \), which is responsible for broadening, is about twice as large in E866 as in E772 data.

Notice that the \( p_T \)-dependence of the ratios in Fig. 3 has the typical shape predicted for the regime of saturation (of quarks), or color glass condensate [29]. Therefore, it might be tempting to interpret such data as an observation of this phenomenon. We warn again that this would be a wrong conclusion, since coherence and saturation are impossible in the short coherence length regime; initial state interaction effects imitate saturation.

Long coherence length

In this regime, \( l_c \gg R_A \), the amplitudes \( \gamma^* \rightarrow l\bar{l} \) for radiating a heavy photon from different nucleons interfere with each other. This interference, which modifies the \( p_T \) distribution of radiation, is known as the Landau-Pomeranchuk effect and is a genuine effect of saturation. Alternatively, one may say that a long-lived fluctuations...
tion $q \rightarrow \gamma^* q$ that propagates through the entire nucleus may be freed as a result of the multiple interactions, in contrast to what happens in the regime of short $l_c$, where the fluctuation is freed after an interaction with a single nucleon. These multiple interactions also lead to a different type of broadening mechanism. Its origin can be understood from the fact that most of the $q \rightarrow \gamma^* q$ fluctuations in the quark are invisible in the sense that their interactions with the target are so weak that the target has restricted resolving power or ability to discriminate a hard $q\gamma^*$ fluctuation, with large intrinsic transverse momentum $k_T$, from a bare quark. However, because of the larger momentum transfer allowed by the multiple interactions in the nuclear medium, the resolution provided by the nucleus is greater than that of a free nucleon. The nucleus can therefore free harder fluctuations than a free nucleon. This is the source of nuclear broadening in the large coherence length regime.

The same effect can be interpreted within the color dipole approach in terms of color filtering. The nucleus filters out projectile dipoles of smaller size than is possible on a free nucleon, leading to larger transverse momenta. In this case, instead of Eq. (7), the transverse momentum distribution of dileptons radiated by a quark interacting with a nucleus is given by the factorized dipole formula [32].

$$
\frac{d\sigma}{dq_T^2} = \int d^2 b \int d^2 r_1 d^2 r_2 e^{i\Phi(q_T)}(r_1-r_2) \times \Psi_{\gamma^* q^*}(r_1, \alpha) \Psi_{q^* q^*}(r_2, \alpha) \left[ 1 - e^{-\frac{1}{2t^2}} \Psi_{N}^a(1,2,x_2) T_A(b) - e^{-\frac{1}{2t^2}} \Psi_{N}^a(1,2,x_2) T_A(b) \right].
$$

(11)

Here $\Psi_{\gamma^* q^*}(r, \alpha)$ is the pQCD calculated light-cone distribution amplitude, which describes the dependence of the $|q\gamma^*\gamma^*\rangle$ Fock component on transverse separation and fractional momentum. The nuclear thickness function at impact parameter $b$ is given by the integral of the nuclear density along the quark trajectory, $T_A(b) = \int dz \rho_A(b, z)$.

The ratio of the nuclear to nucleon cross sections of the Drell-Yan reaction calculated in [32] has a shape typical for the Cronin effect. Such a modification of the transverse momentum dependence is a genuine effect of saturation (or its onset). In spite of the similarity with the analogous results in the short $l_c$ regime, the mechanisms and their formulations in Eqs. (7)-(11) are different. Additionally, the magnitude of the Cronin effect in the long $l_c$ regime is smaller [20].

**ANALYSIS OF E772/866 DATA FOR DRELL-YAN REACTION**

Here we determine the nuclear broadening $\Delta(p_T^2)$ of dileptons produced on different nuclear targets relying on data for $p_T$-dependent ratios depicted in Fig. 3. These include the nucleus-to-deuterium ratios $R_A/D$ from the E772 [2] experiment, and nucleus-to-beryllium ratios $R_A/Be$ from the E866 experiment [31], both at 800 GeV. These results are quite reliable since they are least affected by systematic uncertainties. The ratios $R_i$ are available in 7 bins in $p_T$, $i = 1 - 7$. Correspondingly, one can calculate the broadening, Eq. (3), as

$$
\Delta\langle p_T^2 \rangle = \frac{\sum_i \sigma_i R_i p_T^2 p_T}{\sum_i \sigma_i R_i} - \frac{\sum_i \sigma_i p_T^2}{\sum_i \sigma_i},
$$

(12)

where for each $i$-th bin

$$
\sigma_i = \int d^2 p_T \frac{d\sigma_{\gamma^* q^*}}{d^2 p_T}.
$$

(13)

with $p_T^2$ the mean value of $p_T^2$ within the $i$-th bin. In addition to the seven $p_T$-intervals in data depicted in Fig. 3 we have included an eighth interval in Eq. (12) covering $(p_T^2)^{\text{max}} < p_T < \infty$ and assumed that $R_8 = 1$ in accordance with perturbative QCD expectations [21] (see, however, next section). Since the cross section falls steeply at large $p_T$, this point should not substantially affect the results.

The differential cross section for the Drell-Yan reaction on a free nucleon, $d\sigma_{\gamma^* q^*}/d^2 p_T$, is more affected by systematic uncertainties than the ratio is. We rely on the recently published results of the E866 experiment for the differential cross section in $pp$ and $pD$ collisions [23]. We use this data in [13] for the analyses of all data presented in Fig. 3.

We parametrize the differential cross section of the Drell-Yan reaction as

$$
\frac{d\sigma}{dp_T^2} = N \left( 1 + \frac{p_T^2}{\lambda^2} e^{-p_T^2/\lambda^2} \right)^n.
$$

(14)

The factor in the numerator is introduced to describe a possible forward minimum in the cross section indicated by data [23].

We found that the shape of the $p_T$ dependence does not vary with dilepton mass and $x_2$, as one can see in Fig. 4. A global fit to Drell-Yan data from $pD$ collisions measured by the E866 experiment for different $M$ and $x_2$ bins [23] with common parameters $n$, $\lambda$, and different normalization factors $N$, led to quite a good description, $x^2 = 83.3$ with $N_d.o.f. = 52$. The values of the parameters are $\lambda_1 = 0.78 \pm 0.08$; $\lambda_2 = 0.65 \pm 0.04$; $\lambda_3 = 3.31 \pm 0.21$; $n = 7.00 \pm 0.66$. The results of the fit and the data are depicted in Fig. 4. The error bars include both statistical and systematic errors added in quadrature.

Note that at large $p_T$ the $pD$ cross section in Eq. (14) is falling as $p_T^{-14}$, much steeper than one may expect from
p perturbative QCD. This happens due to restrictions imposed by finiteness of energy, namely, both \( x_1 \) and \( x_2 \) rise with \( p_T \) towards the kinematic limit causing an additional suppression of the cross section \([20]\). Additionally, available data cover a rather restricted interval of \( p_T < 4 \text{ GeV} \).

With the cross section in Eq. (14) we performed an analysis of data from both experiments, E772 and E866. While the ratios measured in E772 are free from major systematic uncertainties, the absolute cross sections are substantially less reliable. Indeed, comparison of the differential cross sections of the Drell-Yan process measured for \( p - D \) collisions in the two experiments reveals substantial differences. Therefore, we rely on the \( p - D \) cross section from the E866 experiment. The results of such a combined analysis of the E772 ratios are depicted in Fig. 5 by closed squares.

The ratios measured in the E866 experiment are for a nucleus relative to beryllium. We use the chain relation in Eq. (15) we performed an analysis of data from both experiments, E772 and E866. While the ratios measured in E772 are free from major systematic uncertainties, the absolute cross sections are substantially less reliable. Indeed, comparison of the differential cross sections of the Drell-Yan process measured for \( p - D \) collisions in the two experiments reveals substantial differences. Therefore, we rely on the \( p - D \) cross section from the E866 experiment. The results of such a combined analysis of the E772 ratios are depicted in Fig. 5 by closed squares.

The ratios measured in the E866 experiment are for a nucleus relative to beryllium. We use the chain relation between the two ratios,

\[
R_{A/D} = R_{A/Be} R_{Be/D} ;
\]

however, the beryllium-to-deuterium ratio is unfortunately lacking. Nevertheless, this cannot be a significant correction, since beryllium is a light nucleus. Indeed, the carbon to deuterium ratio depicted in Fig. 3 is compatible with no nuclear effects at all. If we simply fix \( R_{Be/D} = 1 \), the broadening given in Eq. (12) will be underestimated. However, one can do better by relying on a theoretically predicted value for \( R_{Be/D} \) with Eq. (7). Of course, this introduces an uncertainty related to model dependence, but this uncertainty should not exceed few percent of the observed broadening. The results of this analysis of the E866 data are depicted in Fig. 5 by empty squares.

Notice that the cross section ratios are affected by systematic uncertainties much less than the absolute cross sections. The ratios measured in both experiments have mainly an overall systematic uncertainty 1% in E866 \([34]\), and 2% in E772 \([2]\) measurements. One can see from Eq. (12) that an overall variation of all \( R_i \) leaves the broadening unchanged. Therefore we could disregard the systematic errors of the ratios.

If we fit the data in Fig. 5 assuming a linear dependence on \( A^{1/3} \), in accordance with \([4]\), and do this separately for the E772 and E866 data, we obtain,

\[
D(E772) = (0.029 \pm 0.008) (\text{GeV}/c)^2 ;
\]
\[
D(E866) = (0.059 \pm 0.009) (\text{GeV}/c)^2 .
\]
an analysis [3] depicted in Fig. 1 although with a considerably larger error. However, the E866 data suggest a much stronger broadening. Nevertheless, one may consider the results Eq. (16) as nearly consistent since they lie within two standard deviations. If the broadening according to the E866 data were actually about twice as large as that given by the previous analysis [8] of the E772 data, this would settle the contradiction with theoretical expectations and conclusions drawn from the data for the broadening of J/ψ and Υ mentioned above.

BEYOND THE DATA

Although QCD predicts no nuclear effects at high p_T, the data in Fig. 3 show considerable nuclear enhancement at intermediate values of p_T. This is the well-known Cronin effect first observed for hadronic beams [19]. It is not clear from the data how far in p_T this effect can propagate. The assumption made above for the last unobserved p_T interval, R_8 = 1, is probably incorrect, therefore the broadening effect shown in Fig. 3 could be underestimated. The ad hoc shape of R(p_T) used in [3] is difficult to justify and does not suggest a solution to the problem.

In order to get a hint about how much of the effect might have been missed, one may make use of a theoretically calculated value for R_8. We rely on the dipole formalism, Eqs. (7)-(9), from which one is able to calculate the p_T dependence of the Drell-Yan reaction on proton and nuclear targets in a parameter-free way and explain the data in Fig. 3 well with no adjustment of parameters. The dipole approach also provided the first quantitative explanation of the Cronin effect in available data for hadron-nucleus collisions and correctly predicted the effect for RHIC [21]. Therefore we believe that a calculation of R_8 should be rather reliable. Doing this, Eq. (12) leads to somewhat larger values for the broadening.

The fit of Eq. (4) to this data leads to new values of parameter D,

\[ \bar{D}(E772) = (0.038 \pm 0.008) \text{ (GeV/c)}^2; \]
\[ \bar{D}(E866) = (0.069 \pm 0.009) \text{ (GeV/c)}^2. \]  

This procedure involving theoretical calculations may be considered as a means to estimate the systematic uncertainty. The sign (positive) we obtain for the uncertainty is certainly correct, and its magnitude is not large. Note also that these correlations are corrected for the two sets of data, and therefore one must apply them consistently to both the the results of the E772 and E866 experiments.

BROADENING VERSUS COHERENCE LENGTH

The mechanism of broadening is expected to correlate with the coherence length $\alpha_s [8, 20, 21]$ as described above. The coherence length of the Drell-Yan process is controlled by the value of $x_2$, $l_c \sim 1/m_N x_2 [13]$. Thus, at large $x_2$ the dilepton is produced inside the nucleus and only half of the nuclear thickness contributes to broadening. However, at small $x_2$ the Drell-Yan pair is radiated outside the nucleus and the entire nuclear thickness is effective. It would be a spectacular observation to see a manifestation of this effect in data.

Using recent data from the E866 experiment for the $p_T$-dependent ratios binned in $x_2$ [34, 35] we can also test the predicted $x_2$-independence of broadening. We have performed a similar analysis of the Fe/Be and W/Be ratios in different $x_2$ bins and the results are depicted in Fig. 6.

FIG. 6: Broadening of the mean transverse momentum squared in the Drell-Yan reaction measured [34, 35] on tungsten and iron targets relative to beryllium, $\Delta(p_T^2) = \langle p_T^2 \rangle_A - \langle p_T^2 \rangle_{Be}$.

Unfortunately, within the rather large statistical errors one cannot draw a definite conclusion regarding a possible variation of broadening with $x_2$.

CONCLUSIONS

Analysis of data from the E866 experiment for the Drell-Yan reaction on nuclear targets has resulted in a determination of the magnitude of the nuclear broadening parameter $D(E866)$ given in (16) and (17), which is about twice as large as the one suggested by the E772 data. This result helps to resolve the controversies that were a partial motivation for this analysis. One of these was the expectation that the observed broadening in Drell-Yan would be a factor of 4/3 less than that observed in J/ψ and Υ,
as dictated by the Casimir factor. The updated result is consistent with this expectation and also with the expectation of a rising energy dependence of broadening. And, last but not least, the value of $D(E866)$ agrees with the parameter-free theoretical prediction of \[3\].

This new result also agrees with broadening observed by the HERMES experiment at HERA in hadron production in deep-inelastic scattering on nuclei \[1\]. Indeed, both measurements are well described within the same color-dipole approach \[36\].

Since most of the data from both experiments E772 and E866 correspond to the regime of short coherence length, the observed broadening is directly connected via the relation in Eq. \[9\] to the broadening of a quark propagating through the nuclear medium. In the regime of long coherence length the mechanism of broadening changes and takes the form of color filtering in hard coherent scatterings. To see a possible variation of broadening as function of coherence length, we performed an analysis of data additionally binned in $x_2$. Unfortunately, the restricted statistics did not allow us to draw a definite conclusion from this analysis.

Broadening of the transverse momentum of a parton in a medium leads to induced gluon radiation, i.e., to induced energy loss. This is the basis of the effect called jet quenching which is considered to be a sensitive probe for properties of matter created in heavy ion collisions.

Note that broadening itself is much less dependent on models for hadronization, than jet quenching. Therefore it is a better and more certain probe.

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