Does the A dependence in high-p_T jets come from the European Muon Collaboration effect?

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We show that the models for the European Muon Collaboration effect without a cumulative region ($x > 1$) are not compatible with the "anomalous nuclear enhancement" of high-transverse-momentum jet cross sections. This assertion depends only on the single-hard-scattering picture without final-state interactions or multiscattering and is independent of the exact forms of the nuclear structure functions. We also show that it is possible to reproduce a part of the nuclear enhancement if one includes the cumulative effects.

The "anomalous nuclear enhancement" (ANE) of hadronic jet cross sections[1] has been known for a long time. The experiment involves firing a hadron B at a target nucleus A and looking for hadronic jets with high transverse momentum p_T. The ratio

$$\rho_{\rm jet}(p_T)=\frac{E}{A}\frac{d^3\sigma^{BA}}{dp^3}\Big/ E\frac{d^3\sigma^{BN}}{dp^3} \tag{1}$$

(where E and p are the energy and momentum of the detected jet) of invariant cross sections with a nuclear target and a nucleon N is found to behave as $A^{\alpha(p_T)}$ with $\alpha(p_T)$ greater than 1 and increasing linearly with p_T.

Great excitement has been generated more recently by the discovery of the nontrivial A dependence in deep-inelastic-scattering (DIS) cross sections for leptons on nuclei.[2] This is now called the European Muon Collaboration (EMC) effect after the group that first discovered it. It obviously has repercussions on any high-energy scattering process which used nuclear targets. In particular, the connection between ANE and the EMC effect has been examined by various authors.[3,4] The results obtained up to now seem to be extremely model dependent. We point out here certain model-independent relations and demonstrate that in the absence of multiscattering effects in jet cross sections, the EMC effect and ANE can be shown compatible in the presence of a cumulative region, i.e., by structure functions remaining nonzero at $x > 1$.

Jet cross sections are computed in a hard-scattering picture neglecting all mass scales such as the internal k_T of partons in A and B, quark masses, and Q^2 evolution of the parton densities. Then the invariant jet-production cross section per nucleon is

$$\frac{E}{A}\frac{d^3\sigma^{BA}}{dp^3}=\sum_{ab}\int dx_1dx_2 f_a^A(x_1)f_b^B(x_2)\times\frac{d\sigma_{ab}}{d\hat t}\delta(\hat s+\hat t+\hat u)\ , \tag{2}$$

where x_1 is the fraction of the momentum of a nucleon in A carried by a, x_2 the fraction of the momentum of B carried by b, $f_a^A(x_1)$ the distribution function of *a per nucleon* in the nucleus A, $f_b^B(x_2)$ the density of b in B, $\hat s,\hat t,\hat u$ the Mandelstam variables for the subprocess, and $d\sigma_{ab}/d\hat t$ the subprocess cross section, and the summation is over all species of partons. It is now known[5] that the subprocess cross sections $d\sigma_{ab}/d(\cos\theta)$ (where θ is the scattering angle in the a,b center-of-mass frame), and hence $d\sigma_{ab}/d\hat t$, is to a very good approximation independent of the subprocess (up to numerical factors) and is of the Rutherford form. In fact,

$$\frac{d\sigma_{ab}}{d\hat t}=w_a w_b\frac{d\sigma}{d\hat t}\ ,$$

$$w_q=w_{\bar q}=\tfrac{4}{9},\quad w_g=1\ ,$$

and $d\sigma/d\hat t$ goes as $1/\hat t^2$. Then the sums in (2) can be factored to give

$$\frac{E}{A}\frac{d^3\sigma^{BA}}{dp^3}=\int dx_1dx_2F_A(x_1)F_B(x_2)\frac{d\sigma}{d\hat t}\delta(\hat s+\hat t+\hat u)\ , \tag{3}$$

where $F_i(x)=\sum_a w_a f_a^i(x)$ and $i=A,B,N$. The δ function relates x_1 and x_2 and allows us to integrate over x_2 in (3) giving

$$\frac{E}{A}\frac{d^3\sigma^{BA}}{dp^3}=\int_{x_{1o}}^1 dx_1F_A(x_1)F_B(x_2)\frac{d\sigma}{d\hat t}\ , \tag{4}$$

where $x_2=x_1/(2x_1/x_T-1)$ and the lower limit on the integral $x_{1o}=1/(2/x_T-1)$ comes from putting x_2 to 1. Here $x_T=2p_T/\sqrt{s}$, where $\sqrt{s}$ is the energy available to the beam-nucleon system in its center-of-mass frame and the jet is observed at 90° in the beam-nucleon center-of-mass frame.

We define the ratio of parton distributions in a nucleus and nucleon

$$\tilde\rho(x)=F_A(x)/F_N(x)\ . \tag{5}$$

Although this is not the EMC ratio $\rho_{\rm EMC}$, it becomes indistinguishable from it at $x\gtrsim 0.3$ because the gluon distribution is much softer than that of the quarks.

We can now rewrite (3) as

$$\frac{E}{A}\frac{d^3\sigma^{BA}}{dp^3}=\int_{x_{1o}}^1 dx_1\tilde\rho(x_1)G(x_1,x_T)\ , \tag{6}$$

where

$$G(x_1,x_T)=F_N(x_1)F_B(x_2)\frac{d\sigma}{d\hat t}$$

is non-negative in the range of integration and becomes zero only at the end points. $G(x_{lo},x_T)$ is zero because F_B vanishes when $x_1=x_{lo}$.

The ratio of jet cross sections (1) for A and N is then

$$\rho_{jet}(x_T)=\frac{\int_{x_{lo}}^{1}dx\,\tilde{\rho}(x)G(x,x_T)}{\int_{x_{lo}}^{1}dx\,G(x,x_T)}\ . \tag{7}$$

The variation of $\rho_{jet}(x_T)$ with x_T is obtained by looking at the derivative

$$\frac{d\rho_{jet}}{dx_T}=\frac{1}{\int_{x_{lo}}^{1}dxG(x,x_T)}\left[\int_{x_{lo}}^{1}dx\tilde{G}(x,x_T)\tilde{\rho}(x)-\left(\frac{\int_{x_{lo}}^{1}dx\,\tilde{\rho}(x)G(x,x_T)}{\int_{x_{lo}}^{1}dx\,G(x,x_T)}\right)\int_{x_{lo}}^{1}dx\,\tilde{G}(x,x_T)\right]\ , \tag{8}$$

where $\tilde{G}(x,x_T)=\partial G/\partial x_T$. Putting in the Rutherford form for $d\sigma/d\hat{t}$ we find that

$$\tilde{G}(x,x_T)=G(x,x_T)H(x,x_T) \tag{9a}$$

with

$$H(x_1,x_T)=\frac{2}{x_T{}^2}x_2{}^2\frac{F_B'(x_2)}{F_B(x_2)}-\frac{2}{x_T}\ . \tag{9b}$$

Since structure functions for both protons and pions fall off for $x_2>0.25$, $F_B'(x_2)$ is negative as long as we restrict ourselves to $x_T\gtrsim 0.4$. Then

$$H_+(x,x_T)=-H(x,x_T)=\frac{2}{x_T{}^2}x_2{}^2\frac{|F_B'(x_2)|}{F_B(x_2)}+\frac{2}{x_T}$$

is always positive and increases with x_2 at fixed x_T (Ref. 6) (hence decreasing with x_1). Further, $-\tilde{G}(x,x_T)$ $[=|\tilde{G}(x,x_T)|]$ is non-negative in this region. Defining the mean values

$$\tilde{\rho}(x^*)\int_{x_{lo}}^{1}dx\,G(x,x_T)=\int_{x_{lo}}^{1}dx\,\tilde{\rho}(x)G(x,x_T),\quad \tilde{\rho}(x^{**})\int_{x_{lo}}^{1}dx|\tilde{G}(x,x_T)|=\int_{x_{lo}}^{1}dx\,\tilde{\rho}(x)|\tilde{G}(x,x_T)|\quad (x_{lo}<x^*,\ \ x^{**}<1)\ ,$$

we can rewrite (8) as

$$\frac{d\rho_{jet}}{dx_T}=\frac{\int_{x_{lo}}^{1}dx|\tilde{G}(x,x_T)|}{\int_{x_{lo}}^{1}dxG(x,x_T)}[\tilde{\rho}(x^*)-\tilde{\rho}(x^{**})]\ . \tag{10}$$

For $\rho_{jet}(x_T)$ to increase monotonically with x_T as observed, the term in the square brackets must always be positive. The sign of this term is controlled by the behavior of $\tilde{\rho}$ (which being equal to ρ_{EMC} for $x\gtrsim 0.3$ falls up to $x\simeq 0.7$ and then rises) and the relative positions of x^* and x^{**}. We define

$$M(G;t)=\int_{x_{lo}}^{t}dx\,G(x,x_T)\Big/\int_{x_{lo}}^{1}dx\,G(x,x_T)\ ,$$

$$M(\tilde{G};t)=\int_{x_{lo}}^{t}dx|\tilde{G}(x,x_T)|\Big/\int_{x_{lo}}^{1}dx|\tilde{G}(x,x_T)|\ .$$

It can be easily seen that $M(\tilde{G},t)>M(G,t)$ for all t when $x_T\gtrsim 0.4$ since H_+ is monotonically decreasing with x. This means that $\tilde{G}$ picks up contributions from lower values of x than G and hence

$$x^{**}<x^*\ .$$

Finally, we put in explicit parametrizations of proton and pion structure functions[7] and estimate x^*. It turns out, for example, that $M(G,t)$ for $x_T=0.4$ becomes 0.98 at $t<0.65$, meaning that x^* is between x_{lo} $(=0.25)$ and 0.65. $\tilde{\rho}(x^{**})$ is then greater than $\tilde{\rho}(x^*)$ and hence ρ_{jet} decreases with x_T here. Moreover, going to $x_T=0.8$, we see that both x^* and x^{**} must be greater than 0.67 (x_{lo}) and hence lie in the region where $\tilde{\rho}$ increases. Thus, in this region ρ_{jet} increases with x_T. Such a nonmonotonic behavior of ρ_{jet} is in direct contradiction to experiments.

Thus, the EMC effect is not compatible with the anomalous nuclear enhancement of high-p_T jet cross sections when structure functions are restricted to the region $0\leq x\leq 1$. It should be noted that this conclusion is independent of the gluon distributions in nuclei since our arguments apply to a high-x_T region of scattering where the gluonic contribution to the cross section is small. It also depends on the behavior of $\tilde{\rho}$ (and hence of ρ_{EMC}) only in the region $0.3<x<0.8$, where it is known unambiguously from experiments. Further, the 10–15% effect on structure functions is magnified by the factor of

$$\int_{x_{lo}}^{1}dx|\tilde{G}(x,x_T)|\Big/\int_{x_{lo}}^{1}dx\,G(x,x_T)$$

which turns out to be between 20 and 30 so that ρ_{jet} is much more strongly dependent on x_T.

Next, we investigate the effect of a cumulative region in the quark and gluon densities in a nucleus. Since the nuclear structure functions can now go up to x greater than 1 (Ref. 8) (say x_0), we rewrite (3) as

$$\frac{E}{A}\frac{d^3\sigma^{BA}}{dp^3}=\int_{x_{lo}}^{x_0}dx_1F_A(x_1)F_B(x_2)\frac{d\sigma}{d\hat{t}}\ .$$

Splitting the region of integration into one from x_{lo} to 1 and another from 1 to x_0, and using the definition of $\tilde{\rho}$ in the first region, we can write this as

$$\frac{E}{A}\frac{d^3\sigma}{dp^3}=\int_{x_{lo}}^{1}dx\,\tilde{\rho}(x)G(x,x_T)+\int_{1}^{x_0}dx_1F_A(x_1)F_B(x_2)\frac{d\sigma}{d\hat{t}}\ . \tag{11}$$

The ratio of jet cross sections (1) now becomes

$$\rho_{jet}(x_T)=\rho_{jet}^0+\frac{\int_{1}^{x_0}dx_1F_A(x_1)F_B(x_2)d\sigma/d\hat{t}}{\int_{x_{lo}}^{1}dx\,G(x,x_T)}\ , \tag{12}$$

where ρ_{jet}^0 is the term in ρ_{jet} present in (7). The properties of ρ_{jet}^0 that we have used earlier remain the same even in the presence of the extra term in (12) which we call ρ_{jet}^x and which comes only from the cumulative region of structure functions. The derivative is

$$\frac{d\rho_{\rm jet}}{dx_T}=\frac{d\rho_{\rm jet}^0}{dx_T}+\frac{d\rho_{\rm jet}^x}{dx_T} \quad . \tag{13}$$

$d\rho_{\rm jet}^0/dx_T$ has been computed in (10) and seen to be negative. We now look at the second term. Using mean values as before, we can write

$$\frac{d\rho_{\rm jet}^x}{dx_T}=\rho_{\rm jet}^x(x_T)[H(x^{++},x_T)-H(x^+,x_T)] \ ,$$

where $H(x^{++},x_T)$ is the mean value of $H(x_1,x_T)$ by the weighting function $F_A(x_1)F_B(x_2)d\sigma/d\hat{t}$ and $1\le x^{++}\le x_0$; $H(x^+,x_T)$ similiarly is the mean value by $G(x_1,x_T)$ and $x_{\rm lo}\le x^+\le 1$. The difference term now becomes

$$\frac{2}{{x_T}^2}\left({x_2}^2\frac{F_B'(x_2)}{F_B(x_2)}\bigg|_{x_1=x^{++}} - {x_2}^2\frac{F_B'(x_2)}{F_B(x_2)}\bigg|_{x_1=x^{+}}\right) \quad . \tag{14}$$

Since we are dealing with $x_T>0.4$, $F_B'(x_2)$ evaluated for $x_1=x^+$ is negative. The first term may be either positive or negative. If it is positive then the derivative in (13) is certainly positive. For x^{++} close to 1, the first term may become negative. However, as noted earlier, ${x_2}^2|F_B'(x_2)|/F_B(x_2)$ then decreases with x_1. Hence the expression (14) is always positive, implying that $d\rho_{\rm jet}^x/dx_T$ is positive.

Putting (13) and (10) together, it is seen that the experimentally observed trend of the anomalous nuclear enhancement of jet cross sections (i.e., $d\rho_{\rm jet}/dx_T>0$) can be obtained by simultaneously taking into account the EMC effect at $x<1$ and the cumulative region in nuclear parton densities if

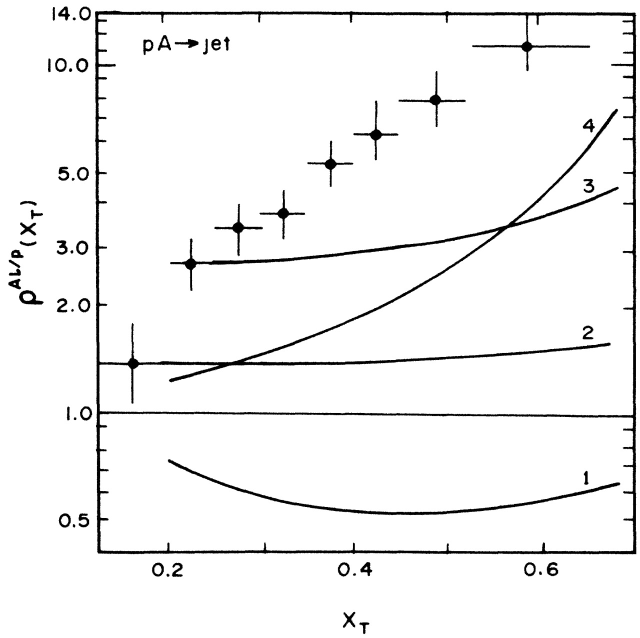

FIG. 1. Computation of $\rho_{\rm jet}^{Al/P}(x_T)$ in various models (1, Ref. 9; 2, Ref. 10; 3, Ref. 11; and 4, Ref. 12) compared to data of Ref. 1. Models 2, 3, and 4 contain a cumulative region and 1 does not.

$$\int_1^{x_0}dx_1F_A(x_1)F_B(x_2)\frac{d\sigma}{d\hat{t}}>\int_{x_{\rm lo}}^1dx\,|\tilde{G}(x,x_T)|\left(\frac{|\tilde{\rho}(x^*)-\tilde{\rho}(x^{**})|}{H(x^{++},x_T)-H(x^+,x_T)}\right) \quad . \tag{15}$$

If the structure function falls off very fast in the cumulative region then the nuclear suppression of cross sections due to the EMC effect at $x<1$ cannot be overcome. The cumulative effect must be strong enough to overcome this suppression. Given a model of nuclear parton densities, we can thus say whether it gives rise to any nuclear enhancement of high-p_T jet cross sections. On the other hand, considerations such as this cannot tell us the exact numerical values of $\alpha(x_T)$. One has to resort to a complete calculation using (2) for such information. Figure 1 shows the results of such a computation[3] and illustrates several points we make here.

In conclusion, we note that in the hard-scattering picture, the EMC effect, and ANE in high-p_T jets are incompatible unless the cumulative region of structure functions in nuclei is taken into account. Even the cumulative effect is not always enough—the structure function at $x>1$ must be strong enough to overcome the effects of nuclear suppression of high-p_T jet cross sections arising from the EMC effect at $x<1$. Thus, this analysis provides a common framework for understanding the results of complete numerical computations done previously.[3,4] These bear out the conclusions we have arrived at here. It should also be noted that a small nuclear effect in structure functions can produce a large x_T dependence of $\alpha(x_T)$. Final-state interactions and multiscattering effects can play an important role in high-p_T production from nuclei. However, a clear theoretical understanding of these effects is lacking. Since we find that structure function effects via a cumulative region can contribute to ANE and these can be measured in deep-inelastic scattering, they constitute a known (in principle) part of ANE. Hence, incorporating this effect would help us to arrive at a better understanding of the other effects mentioned. It is even possible that cluster models for the EMC effect and multiscattering models describe the same physics.

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[6]This is true as long as the logarithmic derivative of $F_B(x_2)$ falls off slower than $1/x_2{}^2$. In particular, it is true for the normal Buras-Gamers parametrization in the form $(1-x_2)^\beta$ for all $\beta > 0$ as well as for $e^{-\lambda x_2}$ for $\lambda > 0$.

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[8]The scaling variable in terms of which perturbative QCD is done is not this x but $\tilde{x} = x/x_0$. The scaling variable $\tilde{x}$ always lies between 0 and 1. The precise value of x_0 depends upon the model but must be less than the nuclear mass A.

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