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Anomalous nuclear enhancement of inclusive spectra at large transverse momentum

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We propose a parton-model interpretation of the "anomalous" nuclear enhancement of inclusive spectra, observed by Cronin *et al.* We argue that the picture representing a nucleus as a collection of quasifree nucleons in slow relative motion is incorrect when the nucleus is probed during a very short time. Our conjecture rests on an extension to nuclei of the parton model of Kuti and Weisskopf. We give a list of observable predictions concerning both hadronic and leptonic interactions with nuclei.

I. INTRODUCTION

There is good evidence that hadronic forces have a finite range in rapidity. This does not mean, of course, that two energetic hadrons do not interact, since the colliding particles dissociate virtually and those among their respective constituents which are separated by a small distance in rapidity feel the forces.[1] A corollary[2]: The nuclear matter is transparent for hard partons, and therefore the distribution of fast debris of a hadronic projectile should be the same in an elementary and in a nuclear collision; this prediction is corroborated by data.[3] The argument holds for rapidities larger than a certain critical value $y_c \sim \ln A$, where A is the nuclear mass number, if one neglects the constraints due to unitarity. The interaction of hard partons, neglected in the first approximation, is believed to be responsible for production of secondaries at large transverse momentum. As long as the nucleus can be regarded as a collection of instantaneously free nucleons in slow relative motion, one predicts that at large transverse momentum k_T

$$\frac{\omega d\sigma}{d^3k}(NA\rightarrow hX)\approx A\,\frac{\omega d\sigma}{d^3k}(NN\rightarrow hX)\,. \qquad (1)$$

(Since no confusion is possible, in order to simplify the notation the letter A is occasionally used in this paper to represent a nucleus with mass number A; a nucleon is represented by the letter N.) The above introduction is purposely sketchy, since we assume that the reader is familiar with the parton-model lore. When a more precise formulation of the parton model will be needed we shall adopt the constituent-interchange model (CIM) of Blankenbecler, Brodsky, and Gunion.[4]

Actually, it is observed[5] that

$$\frac{\omega d\sigma}{d^3k}(NA\rightarrow hX)> A\,\frac{\omega d\sigma}{d^3k}(NN\rightarrow hX) \qquad (2)$$

for k_T large enough. For example, the left-hand side of (2) is larger than the right-hand side by almost a factor of 2, for pion production in proton collisions with a tungsten target at 300 GeV and for $k_T > 3.5$ GeV/c. This phenomenon will be referred to as the *anomalous nuclear enhancement* (ANE). Several authors[6] called upon various multiple-scattering effects in order to explain this phenomenon. In this paper we shall argue that the picture representing a nucleus as a collection of quasifree nucleons is wrong when the nucleus is probed during a very short time and that this is the true origin of ANE. Basic to our argument is the conjecture that there are more energetic "sea" constituents in a nucleus than might be naively thought. This conjecture is abstracted from the parton model of Kuti and Weisskopf,[7] extrapolated to the nuclear case, and is to some extent supported by the data on particle production in heavy-ion collisions.[8] The same conjecture also implies a substantial enhancement of the yield of $\mu^+\mu^-$ pairs produced on a nuclear target by the Drell-Yan mechanism $q\bar{q}\rightarrow\mu^+\mu^-$. Such an enhancement has been introduced *ad hoc* to fit the data.[2,9]

II. EXTRAPOLATION OF THE KUTI-WEISSKOPF MODEL

The production of secondaries with large transverse momenta is sensitive to very-short-time effects in nuclear matter. Conventional nuclear physics does not teach us much about these effects and we have to look for hints elsewhere. In the absence of an appropriate theory we follow a conservative approach: We extrapolate models and ideas which turned out to be useful in the study of more elementary particles to the nuclear case.

Let us remark first that, because of the time dilation, the natural frame for discussing short-time effects within a composite system is the frame where the system in question has a very large momentum P. But in such a frame it is appropriate to think of a nucleus as made up of quarks, antiquarks, gluons, etc.

How do the one-parton distribution functions depend upon the nuclear mass number A? We shall seek an answer to this question in the framework of an explicit parton model due to Kuti and Weisskopf. Following these authors, we assume that partons can be regarded as independent particles. By properly choosing their individual wave functions one hopes to mimic the dynamics of the composite system.

Consider first an unspecified hadronic system with M valence quarks and an indefinite number of sea partons. In the simplest case, where one does not distinguish between different types of sea and valence partons, the probability of a configuration including exactly N sea partons is

$$dp_N(x_1, \ldots, x_{M+N}) = Z\,\frac{a^N}{N!}\,\delta\left(1-\sum_{i=1}^{M+N} x_i\right)\left[\prod_{j=1}^{M}\frac{v(x_j)dx_j}{(x_j^2+m_v^2/P^2)^{1/2}}\right]\left[\prod_{j=M+1}^{M+N}\frac{s(x_j)dx_j}{(x_j^2+m_s^2/P^2)^{1/2}}\right], \tag{3}$$

where x_j is the momentum of the jth constituent, in units of P, and the functions $s(x)$ and $v(x)$ have the following behavior near $x=0$:

$$s(0)=1, \quad v(x)\sim x^b \quad (b>0)\,. \tag{4}$$

The physical quantities of interest are calculated in the limit $P\to\infty$ and turn out to be independent of the effective mass parameters $m_{s,v}$. (However, a systematic use in the following of the asymptotic formulas is, strictly speaking, legitimate when $P \gg Am_{s,v}$.) The normalization factor Z is unambiguously determined from self-consistency requirements. The "transverse" motion of partons is ignored for simplicity. A generalization to the case where there are different types of sea and/or valence partons is straightforward.

Consider now a nucleus with nuclear mass number A. Obviously, the number of valence quarks $M=3A$. As in all statistical models, the parameter a is proportional to the volume of the system and, consequently,

$$a=\text{const}\times A\,. \tag{5}$$

The average multiplicity of sea particles is $\sim a\ln(P/m_s)$. Equation (5) means that there is, on the average, A times more constituents in a nucleus than in a nucleon. Finally we assume that the shapes of the individual parton wave functions, as functions of the scaled momentum x, are independent of A. The motivation for this last assumption, certainly in the spirit of the statistical model we are working with, is mostly the economy of thought.

The electromagnetic form factor is proportional to the probability, which is nonzero in the model, that a single quark picks almost the whole momentum of the hadron.[1] On the other hand, a nonzero form factor corresponds to a coherent response of the nucleus. Thus, this coherent response is represented in the model by a dramatic fluctuation in the partition among partons of the total available momentum. Notice that such a fluctuation is incompatible with a simple picture representing a fast-moving nucleus as a collection of A almost-free nucleons with momentum $\approx P/A$ each. However, the simple picture is not necessarily implied by the results of low-energy nuclear physics. The transformation of a wave function from one relativistic frame to another is a nontrivial problem whose solution requires a complete knowledge of the underlying dynamics. Anyhow, the simple picture is in variance with the data: The breakup of heavy ions is associated with relatively frequent production of pions carrying momentum much larger than P/A.[8] We are therefore tempted to speculate that the existence of parton configurations, excluded by the naive picture, represents short-time collective nuclear effects.

III. A TOY MODEL

There is experimental evidence that sea-parton distribution functions $G_{h/N}(x)$ fall much more rapidly toward $x=1$ than the analogous distribution functions corresponding to the valence quarks. [$G_{h/A}(x)dx$ is the average number of partons h, with scaled momentum within the interval $(x, x+dx)$, in a nucleus with mass number A. For $A=1$ we write $G_{h/N}(x)$.] The particularly rapid fall of $G_{h/N}(x)$, when h is a sea constituent, will be essential for our argument. A rationale for this behavior is provided by the so-called dimensional counting rules.[10] We shall often use these rules in the following, in order to constrain the discussion (it is an exploratory work and we do not have to worry whether these rules are exact, or only approximate as some recent data seem to indicate). The counting rules are derived using arguments from quantum field theory. In general, these rules are not satisfied by models of the Kuti-Weisskopf type, which are essentially phase-space models and do not involve constraints characteristic for a field theory. However, one can try to enforce the satisfaction of the counting rules by properly choosing the parameters, which are otherwise almost arbitrary, of a phase-space model. The original model of Ref. 7 cannot reproduce the very different behavior near $x=1$ of valence and sea parton distribution functions. Probably the simplest phase-space model of the nucleon which can be made compatible with counting rules is the one where the nucleon is regarded as built up of three valence quarks and of a sea of virtual "pions" and gluons. We neglect complications due to spin and isospin and put

$$s(x)=\begin{cases}(1-x)\sin^2\eta\\(1-x)\cos^2\eta\end{cases} \tag{6a}$$

for "pions" and gluons, respectively, and

$$v(x)=x^b,\quad b>0 \tag{6b}$$

for quarks. With the help of standard Laplace-transform techniques, one obtains from an analog of Eq. (3), in the limit $P\to\infty$ and for $x>0$,

$$G_{q/N}(x)=\frac{3a^{b/2}}{\Gamma(b)J_{a+3b-1}(2\sqrt{a})}x^{b-1}\times(1-x)^{(a+2b-1)/2}J_{a+2b-1}(2[a(1-x)]^{1/2}), \tag{7a}$$

$$G_{M/N}(x)=\frac{a\sin^2\eta}{J_{a+3b-1}(2\sqrt{a})}x^{-1}\times(1-x)^{(a+3b+1)/2}J_{a+3b-1}(2[a(1-x)]^{1/2}). \tag{7b}$$

We wish to impose the behavior near $x=1$ predicted by the counting rules, viz.

$$G_{q/N}(x)\propto(1-x)^3,\quad G_{M/N}(x)\propto(1-x)^5. \tag{8}$$

We are thus led to set $a=2$ and $b=1$.

We explained in the preceding section how to extrapolate the model to the nuclear case. Hence, we assume that the number of quarks is $3A$ and we replace a by Aa. A calculation analogous to the one which leads to Eqs. (7) yields now, for $a=2A$ and $b=1$, the following result:

$$G_{q/A}(x)=\frac{3\sqrt{2}\,A^{3/2}}{J_{5A-1}(2\sqrt{2A})}\times(1-x)^{(5A-2)/2}J_{5A-2}(2[2A(1-x)]^{1/2}), \tag{9a}$$

$$G_{M/A}(x)=\frac{2A\sin^2\eta}{J_{5A-1}(2\sqrt{2A})}x^{-1}\times(1-x)^{(5A+1)/2}J_{5A-1}(2[2A(1-x)]^{1/2}). \tag{9b}$$

In a realistic theory $G_{h/A}(x)$ should be proportional to $A^{2/3}$ for wee x's. This condition is not satisfied by Eqs. (9). As expected, the model is manifestly incorrect for those values of x where shadowing becomes important. This difficulty can be ignored, however, as long as one is interested in hard-parton scattering only, as we are. We imposed the satisfaction of the dimensional counting rules in the case $A=1$. It is easily found from Eqs. (9) that these rules are no longer satisfied by the toy model when $A>1$.

It is a simple matter to calculate the contribution of the subprocess $qM\to qM$ to the inclusive "pion" production at large k_T, using the CIM formula

$$\frac{\omega d\sigma}{d^3k}(NA\to MX;stu)=\frac{1}{\pi}\int_0^1\int_0^1 dx\,dy\,G_{q/A}(x)G_{M/N}(y)\frac{d\sigma}{dt'}(qM\to qM;s'=Axys,t'=yt,u'=Axu)\,s'\delta(s'+t'+u')$$
$$+\frac{1}{\pi}\int_0^1\int_0^1 dx\,dy\,G_{M/A}(x)G_{q/N}(y)\frac{d\sigma}{dt'}(qM\to qM;s'=Axys,t'=Axu,u'=yt)s'\delta(s'+t'+u'), \tag{10}$$

where s, t, and u are Mandelstam kinematic variables calculated as if the target were a nucleon carrying the fraction $1/A$ of the momentum of the nucleus. (This common choice of kinematic variables is purely a matter of convention.) Dimensional counting implies that

$$\frac{d\sigma}{dt}(qM\to qM;stu)=s^{-4}f(t/s),$$

where the function $f(w)$ depends on the details of the dynamics. In our calculations we set $f(w)=\text{const}$ for definiteness. The dependence of the calculated ANE on the choice of $f(w)$ is essentially trivial. The results of a numerical calculation for the subprocess $qM\to qM$ and of an analogous calculation for $MM\to MM$ are given in the Table I. [As will be explained later on, ANE results from an "anomalous" enhancement of the number of energetic sea partons in a nucleus. In the NN rest frame, the center of mass of the two colliding constituents has the tendency to move in the direction of motion of the nucleus. Once this is realized, it is easy to understand the dependence of ANE on the choice of $f(w)$. Thus, pion

TABLE I. The ratio $(\omega d\sigma/d_3k)(NW \to MX)/[184(\omega d\sigma/d_3k)(NN \to MX)]$ calculated for 300-GeV nucleons incident on tungsten ($A = 184$) and for mesons emitted at 90° in the nucleon-nucleon rest frame.

$x_T = 2k_T/\sqrt{s}$	0.3	0.4	0.5	0.6
The anomalous nuclear enhancement in the toy model.				
Only subprocess $qM \to qM$ taken into account	0.97	1.13	1.57	2.85
Only subprocess $MM \to MM$ taken into account	1.25	1.70	2.84	6.53
The anomalous nuclear enhancement calculated using Eqs. (13), (14), and (16).				
Only subprocess $qM \to qM$ taken into account	0.97	1.05	1.34	2.20
Only subprocess $q\bar{q} \to M\bar{M}$ taken into account	1.44	1.99	3.49	8.72

production at 90° in the NN rest frame is due preferentially to the backward pion-quark scattering, and ANE is strengthened when the relative importance of this backward scattering is increased. For example, with $d\sigma/dt \propto s^{-4} + (su)^{-2}$ the first line of Table I becomes 1.04, 1.24, 1.77, and 3.36.]

We do not claim that our toy model is realistic and we do not attach much importance to the numbers shown in the Table I, except for qualitative trends. The purpose of the model is to serve as an example. Indeed, the toy model predicts a significant anomalous nuclear enhancement of inclusive cross sections at large transverse momentum. Let us examine now in more detail the mechanism of this enhancement.

A closer scrutiny of our numerical computations reveals that the rather complicated expressions (9) conceal a basically very simple behavior. Thus, $G_{M/A}(x)$ behaves very much like $Ax^{-1}(1-x)^{5A}$ for $x \lesssim 1/A$ and essentially for all A. Furthermore, the nuclear parton distribution functions contribute significantly to the integral in (10) when their arguments are $\lesssim 1/A$. This suggests we use, instead of x, the more convenient variable $z = Ax$. Thus

$$G_{M/A}(x)dx \approx \text{const}\times Az^{-1}(1-z/A)^{5A}dz, \quad z = Ax \lesssim 1\,. \tag{11a}$$

For $A = 1$, the above expression reduces to

$$G_{M/N}(z) \approx \text{const}\times z^{-1}(1-z)^5\,, \tag{11b}$$

a form of pion distribution function often used in CIM calculations. However, if a nucleus were simply a box of quasifree nucleons in slow relative motion, (11b) would imply that

$$G_{M/A}(x)dx \approx \begin{cases} \text{const}\times Az^{-1}(1-z)^5 dz, & z \leqslant 1 \\ 0 \text{ otherwise}, \end{cases} \tag{11c}$$

instead of (11a). On the other hand, when $A \to \infty$

$$(1-z/A)^{5A} \to \exp(-5z)\,,$$

and it is a numerical fact that $\exp(-5z)$ falls slower than $(1-z)^5$, especially when z approaches unity. Thus in the toy model *the population of hard virtual pions within a nucleus is enriched compared to a nucleon, by a factor larger than A.* The effect is sufficient to produce a considerable enhancement of production at large k_T, although the average fraction of the total momentum carried by virtual pions changes little (by less than 13%, when A increases from 1 to ∞).

The threshold factor $(1-z)^3$ which dominates the structure of the valence quark distribution function $G_{q/N}(z)$ is also converted into $\exp(-5z)$ when $A \to \infty$. In this case, however, one has an enhancement only when z is close to unity. Otherwise one has a suppression. This is why ANE is much stronger for the subprocess $MM \to MM$ than for $qM \to qM$ (cf. Table I).

IV. BEYOND THE TOY MODEL

A. Behavior of parton distribution functions

We believe that there is no point in constructing more complicated phase-space models of a nucleus in the $P = \infty$ frame. It is more instructive to abstract, from our experience with the simplest of such models, those features which presumably characterize the whole class:

(1) The volume of the relativistic phase space with a cutoff on transverse momenta behaves, modulo logarithmic factors, like a power of the total available momentum, viz. P^{const}. A distribution function $G_{h/A}(x = k/P)$ is proportional to the ratio of phase-space volumes $(P-k)^{\text{const}}/P^{\text{const}} \equiv (1-x)^{\text{const}}$. The exponent contains a term proportional to the volume of the system and another one proportional to the number of valence quarks. It is therefore linear in A.

(2) $G_{h/A}(x)dx$ is also proportional to the product of the modulus squared of the individual wave function of h by the relativistic phase-space element dx/x. This factor determines the behavior of the distribution function for small values of x and can be represented by some power of x times dx.

(3) The average number of constituents of a given type is obtained upon integrating the corresponding distribution function and is proportional to A. This condition constrains the over-all normalization of $G_{h/A}$.

Thus, in the limit $P \to \infty$, the shape of parton distribution functions in statistical models of the Kuti-Weisskopf type is very roughly

$$G_{h/A}(x)dx \approx Ac_h z^{p_h}(1 - z/A)^{r_h}dz\,, \quad z = Ax > O(1/P)\,, \tag{12}$$

or a linear combination of such terms, with c_h weakly dependent on A, p_h independent of A, and r_h linear in A. Remember the notation: z is the momentum of h measured in units of P/A, which is the average momentum of a *nucleon* (for $A = 1$, $z \equiv x$).

We shall assume in the following that the dimensional counting rules also hold for nuclei. A sufficiently elaborated statistical model can satisfy these rules for any A. Then, for $A = 1$, the exponent r_h depends strongly on the constituent's type: $r_h = 3, 5,$ or 7 for h = quark, pion, or antiquark, respectively. On the other hand, for $A \gg 1$, the ratio r_h/A becomes independent of h. In the toy model we had $r_h/A \to 5$. The counting rules imply that $r_h/A \to 6$. The independence of r_h/A of h, as A becomes large, means that *all nuclear parton distribution functions fall, with increasing z, at roughly the same rate* (provided $A \gg 1$).

With the above rules of thumb, the simplest form of the "pion" distribution function $G_{M/N}(z) = c_M z^{-1}(1 - z)^5$ is converted into

$$G_{M/A}(x)dx = Ac_M z^{-1}(1 - z/A)^{6A-1}dz\,, \quad z = Ax > O(1/P)\,, \tag{13}$$

and not into (11c) as one might expect naively. Similarly, the simplest distribution of antiquarks $G_{\bar{q}/N}(z) = c_{\bar{q}} z^{-1}(1 - z)^7$ is converted into

$$G_{\bar{q}/A}(x)dx = Ac_{\bar{q}} z^{-1}(1 - z/A)^{6A+1}dz\,, \quad z = Ax > O(1/P)\,, \tag{14}$$

and not into

$$G_{\bar{q}/A}(x)dx = \begin{cases} Ac_{\bar{q}} z^{-1}(1 - z)^7 dz, & z \leq 1 \\ 0 \text{ otherwise.} \end{cases} \tag{15}$$

The enhancement of the right-hand side of (13) over the right-hand side of (11c), or of (14) over (15), is analogous to that discussed in the context of the toy model.

The case of the quark distribution function is more subtle, since this function involves several contributions behaving differently. In general, the valence-quark contribution is suppressed for large A, while the sea and bremsstrahlung contributions[11] are enhanced. The difficulty in analyzing $G_{q/N}$ is in finding a roughly realistic expression, still simple enough to apply our "rules of thumb." For N = 50% of proton plus 50% of neutron, the information from deep-inelastic lepton scattering is compatible with

$$G_{q/N}(z) = 0.3z^{-1}(1 - z)^7 + 2.1z^{-1/2}(1 - z)^7 + 6.8(1 - z)^3 \tag{16a}$$

and with the assumption that the "sea" is SU(3)-symmetric (see Ref. 11, where a theoretical justification of the different threshold factors is given). This corresponds to the scaled momentum of all nonstrange quarks and antiquarks ≈ 0.49 (experimentally: 0.49 ± 0.07).[12] The ratio of momentum of antiquarks to the total momentum of quarks and antiquarks is ≈ 0.1 (experimentally: 0.1 ± 0.03).[12] The last two terms in (16a) integrate to 3. Finally, the behavior of $\frac{1}{2}\nu(W_2^p + W_2^n)$ is rather well reproduced. The right-hand side of the Eq. (16a) is a linear combination of terms of the form given by (12). Assuming, as before, that the number of sea quarks is proportional to A, that the number of valence quarks is $3A$, and that the momentum carried by valence quarks is independent of A (this dependence is weak in statistical models), one converts (16a) into

$$\begin{aligned} G_{q/A}(x)dx = A\{&0.3z^{-1}(1 - z/A)^{6A+1} \\ &+ [1.4 + O(1/A)]z^{-1/2}(1 - z/A)^{6A+1} \\ &+ [12 + O(1/A)](1 - z/A)^{6A-3}\}\, dz\,, \\ z = Ax > O(1/P)\,. \end{aligned} \tag{16b}$$

Comparing (16a) with (16b) one finds that *the anomalous suppression of the valence-quark contribution and the anomalous enhancement of the sea and bremsstrahlung contributions compensate approximately:* The expression in curly brackets on the right-hand side of (16b) differs little from the right-hand side of (16a) (except for $z \gtrsim 0.8$, where it is larger; all contributions to $G_{q/A}$ are enhanced). This finding is compatible with the data on the A dependence of the deep-inelastic structure functions.[13] However, the experimenters were focusing their attention on the region of small and moderate z in a search for shadowing effects, which were ignored in our discussion. It would be interesting to observe the dependence on A of cross sections for the deep-inelastic lepton scattering at large z, say at $z > 0.8$, where we except an anomalous enhancement.

Finally, the number of hard biquarks in a nucleus is increased, compared to a nucleon, by a factor smaller than A. Indeed, $G_{qq/N}(z)$ falls particularly slowly toward $z = 1$, like $(1 - z)$.

B. A physical picture

Basic to our speculations is the extension of the parton model to nuclei. If one made a snapshot of a nucleus one could not distinguish between virtual particles emitted and eventually reabsorbed by the same nucleon and those exchanged between different nucleons. Therefore, we treat a nucleus as one bag of partons. We believe that there are large fluctuations in the partition among these partons of the total available momentum. As a consequence, those species of partons which have a very small probability of carrying a considerable fraction of the momentum of a nucleon, when this nucleon is free, have their momentum distribution enhanced in a nucleus. This enhancement results, in turn, in an enhancement of cross sections for production of secondaries with large transverse momenta.

Of course, the momentum distributions of all partons cannot be enhanced. In the toy model, the average momentum of sea partons is increased slightly, at the expense of the valence quarks. In the preceding section we assume instead that the increase of the average momentum of sea quarks and antiquarks is realized at the expense of gluons. Anyhow, sea quarks and antiquarks carry a small fraction of the total momentum. This leaves a considerable latitude in the choice of a model.

In the case of a nucleon-nucleon collision, the contribution of a given constituent-interchange subprocess to the production at large k_T depends on the degree of "forbiddeness" of the subprocess (criteria for "forbiddeness" are excellently reviewed in Ref. 4). It is obvious that a subprocess is suppressed if it involves initially a constituent whose distribution function is required to fall rapidly toward $z=1$. For example, if the corresponding amplitudes were of the same magnitude, subprocesses of the type $q\bar{q}\to$ something would be suppressed compared to subprocesses of the type $qM\to$ something, because it is easier to find in a nucleon a hard nonstrange meson [threshold factor $(1-z)^5$] than a hard antiquark [threshold factor $(1-z)^7$]. According to the discussion of the preceding section, all *nuclear* parton distribution functions fall with z at about the same rate (when $A\gg 1$). Therefore, *in nuclear production at large k_T, "forbidden" subprocesses are enhanced relative to "allowed" subprocesses.* The more a subprocess is "forbidden," the more it is enhanced. The discussion of the behavior of different parton distribution functions, presented in the preceding section, also implies that *only those subprocesses which involve at least one sea particle can contribute to the observed anomalous nuclear enhancement.*

Subprocesses involving, in the initial state, two sea partons, such as $MM\to MM$, exhibit a strong ANE. However, the contribution to the inclusive cross section of a single such process is smaller by an order of magnitude, at least, than a contribution of a typical subprocess involving in the initial state one sea constituent and one quark. This is related to the small number of hard antiquarks in a nucleon. We cannot exclude the possibility that the observed ANE is a result of a complicated balance between a myriad of competing subprocesses. If one seeks simplicity, however, one is led to conjecture that subprocesses involving in the initial state valence constituents only, like $qq\to B\bar{q}$ and especially like $q(qq)\to BM$, are unimportant at $\lesssim 500$ GeV. [The subprocess $q(qq)\to BM$ is expected to exhibit an anomalous nuclear suppression.] Once this conjecture is made, one predicts that ANE is particularly strong for the production of K^-, p and $\bar{p}$ in agreement with observation. The subprocess which is favored by our considerations is the "fusion" process $q\bar{q}\to h\bar{h}$, whose importance has already been emphasized in a different context by Landshoff and Polkinghorne.[14] The result of a tentative CIM calculation, using Eqs. (13), (14), and (16) and assuming that $d\sigma/dt=\text{const}\times s^{-4}$, is given in the Table I. Obviously, we also expect that the production of lepton pairs by the Drell-Yan mechanism $q\bar{q}\to\mu\bar{\mu}$ and, in fact, by any other mechanism involving sea partons, is enhanced when the target is a nucleus.

It is rather evident that harder partons are necessary to produce secondaries with larger $x_T=2k_T/\sqrt{s}$. Hence, the relative contribution to a CIM integral of the region, where parton distribution functions exhibit an important nuclear enhancement, increases together with x_T; *in enhanced subprocesses,* ANE *is stronger for larger* x_T. This trend is clearly seen in the Table I. We cannot exclude the possibility that different subprocesses contribute contrariwise to observed cross sections, so as to make the dependence of ANE on x_T complicated. However, an increase of ANE when one moves toward the kinematic boundary is, from our point of view, a very natural prediction. And in reactions where ANE is strongest, this increase of ANE with x_T should be the most evident. In fact, the data of Ref. 5 behave like that. This is manifest when one examines the variation with k_T of the ratio of inclusive cross sections for production on tungsten and berilium respectively, using the data tables given in Ref. 5. [The authors of Ref. 5 fit their data assuming that inclusive cross sections are proportional to $A^{\alpha(k_T)}$. The value of the fitted exponent $\alpha(k_T)$ is only an indirect measure of ANE and its variation with k_T does not represent very accurately the trend of the data.] The increase of ANE is systematic (also for pion production), except for data at the largest transverse momentum, $k_T = 6.1$ GeV/c. At $k_T = 6.1$ GeV/c, ANE is smaller than at 5.34 GeV/c but, because of errors, one can have some doubts about the significance of the effect.

We predict also that the nuclear enhancement of the production of lepton pairs via the Drell-Yan mechanism should be stronger for the larger invariant mass $\sqrt{Q^2}$ of the pair, since increasing Q^2 one again moves toward the kinematic boundary. This is a welcome conclusion, since the discrepancy between data on production of muon pairs on uranium[15] and parton-model calculations[2,9] increases with Q^2. Notice that people who interpret ANE as a rescattering effect have nothing to say about the production of lepton pairs.

It is worth mentioning that, as long as ANE is ignored, the picture of high-energy nuclear interactions of hadrons adopted in this paper is identical to that postulated in the paper by Kühn[6]. However, Kühn interprets ANE as a result of multiple scattering of hard partons in nuclear matter. The effect he discusses can coexist with the nuclear effects examined in this paper. We suspect, however, that an extension of Kühn's calculation to the case of production of heavy particles would make evident a strong dependence of the result on the input information and would be therefore rather artificial, while experimentally the nuclear enhancement effect is strongest for K^-, $\bar{p}$, and p production.

In this paper, we use statistical models of the Kuti-Weisskopf type because we have nothing better at our disposal. We checked that ANE of right magnitude can easily be obtained (cf. Table I). However, the experimentally observed dependence of ANE on A, of the form $A^{\alpha(k_T)}$, cannot be reproduced by models of the Kuti-Weisskopf type, where the structure of parton distribution functions is dominated by the phase-space factor $(1 - z/A)^{\text{const} \times A}$. This failure is due to the rather rapid convergence of $(1 - z/A)^{\text{const} \times A}$ toward $\exp(-\text{const} \times z)$ when A increases, and is a consequence of the extreme crudeness of statistical models. Anyhow, the really intriguing aspect of nuclear production at large k_T is the fact that cross sections are not just proportional to the number of scatterers, as in (1), although one probes distances much smaller than the size of a nucleon (and one has to remember the quasi-absence of cascading within nuclei as well as the rapid fall with k_T of production cross sections). To this phenomenon we propose an interpretation which presents certain advantages over interpretations via multiple scattering and which leads to characteristic observable predictions (see Sec. V). Unfortunately, a successful description of the variation with A of the effect requires a more realistic picture of the nucleus than the one we have used.

V. OBSERVABLE PREDICTIONS

We are aware of the speculative character of this paper. This is one reason why we limit ourselves to a qualitative discussion (remember, nevertheless, that nobody has succeeded as yet in producing a successful fit to the NN data at CERN ISR *and* at Fermilab energies with a CIM model; under these conditions it does not make sense to attempt a fit to nuclear data). However, our considerations lead to predictions, listed below, which can be tested experimentally:

(i) As was already mentioned, in processes exhibiting ANE and in the NN rest frame, the center of mass of the two colliding constituents should have a tendency to move in the direction of motion of the nucleus. Therefore, we expect that any jet structure, which might be observed in pp collisions on the side opposite to the large-k_T particle, should be boosted in the direction of motion of the nucleus.

(ii) Of course, we predict that the production of lepton pairs exhibits ANE in close analogy to the large-k_T production of hadrons.

(iii) Our preference for, or prejudice against, certain constituent-interchange subprocesses has immediate consequences for quantum-number correlations in large k_T production. For example, we would be surprised if the transverse momentum of a jet carrying baryonic number $B=0$ would often be compensated, on the other side, by the transverse momentum of a jet carrying $B=1$. Our considerations favor the fusion processes $q\bar{q} \to h\bar{h}$. Therefore, a jet carrying $B=-1$ should be accompanied by another jet on the other side and carrying $B=1$, etc. (see Ref. 16 for a discussion of quantum-number correlations in CIM).

(iv) We expect an anomalous nuclear enhancement of cross sections for the deep-inelastic lepton scattering at small values of the Bjorken variable $\omega = 1/z$ (see Sec. IV A).

(v) As stated earlier, the more a process is "forbidden"[4] the more it should be enhanced. This general conclusion is supported by the data of Ref. 5, but a more systematic check of this prediction would be interesting.

The importance of studying interactions of very energetic particles with nuclei has been repeatedly emphasized by many people. In particular, these interactions provide invaluable information on the time scale relevant for multiple production processes. Our speculations suggest that they might also reveal unconventional nuclear effects. That such effects exist is already evidenced by the surprising results obtained with heavy ions.

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