QCD EFFECTS IN PARTICLE PRODUCTION ON NUCLEI*

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ABSTRACT

I discuss some questions related to hard scattering processes in nuclei and corrections to the leading twist approximation. The QCD factorization theorem requires that high energy partons do not lose energy while traversing the nucleus. I explain the physical reason for this. The theorem also states that spectator partons, not involved in the hard collision, have no influence on the inclusive cross section. Important spectator effects are, however, seen in the data for certain reactions and in some kinematical regions. I discuss the reasons from a phenomenological and theoretical point of view. Finally, I mention some methods for analyzing hard processes in regions of very large $x$, where the leading twist terms are not dominant.

Studies of the dependence of hard collisions on the nuclear number of the target and/or projectile have revealed a large number of interesting, diverse and often surprising QCD effects. The nucleus has been, and will continue to be, a versatile tool for uncovering reaction mechanisms and indicating where our standard methods are inadequate. In this talk I shall confine myself to two particular questions, related to the energy loss of partons in nuclei and the importance of higher twist corrections due to spectator parton interactions. For a wider view of nuclear effects, I refer to recent reviews\(^1\).

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1. Parton Energy Loss in Nuclei

At leading twist, \textit{i.e.}, up to corrections of \( \mathcal{O}(1/Q^2) \) in the hard scale \( Q \), the dependence on hard and soft physics factorizes in a simple way in QCD\(^2\). Thus, \textit{e.g.}, the cross-section for the inclusive reaction \( a + b \rightarrow c + X \) is written

\[
\sigma(a + b \rightarrow c + X) = \sum_{ijk} F_{i/a}(x_1, Q^2) F_{j/b}(x_2, Q^2) \hat{\sigma}(ij \rightarrow k) D_{c/k}(z, Q^2)
\]

for large values \( Q \) of the transverse momentum and/or mass of the final hadron \( c \). The factorization of the cross section into a product of soft single parton structure functions \( (F) \), hadronization functions \( (D) \) and a hard parton-level cross-section \( (\hat{\sigma}) \), is a remarkable simplification. The reaction rate depends only on the probability of finding a single parton in each colliding particle, and the fragmentation of the final parton is independent of the nature of the target, projectile and the hard subprocess.

A particular consequence of the leading twist Eq. (1) is that the projectile structure function \( F_{i/a}(x_1) \) is independent of the nature of the target \( b \). At first this appears surprising, since one might expect the projectile parton \( i \) to suffer energy loss from gluon radiation while penetrating a large nuclear target \( b \), en route to its hard interaction. However, the nucleus-induced energy loss is suppressed at high energies\(^3\).

To see the physical reason for the low energy loss, consider a typical Fock state of the projectile hadron wave function. It consists of the parton that is going to have the hard collision and a number of spectator partons. Now at high energies, essentially due to time dilation, the various partons in the Fock state generally do not have time to communicate (interact) with one another while in the target — the Fock states do not mix. The mixing time (or “life-time”) of a Fock state is given by the inverse of the energy difference between the mixing states. Thus, the emission of a gluon from a quark, \( q \rightarrow q + G \), involves a kinetic energy difference \( \Delta E \) between the \( q \) and the \( qG \) states given at high energy by

\[
\Delta E(q \rightarrow qG) \simeq \frac{1}{2p} \left[ m_q^2 - \frac{m_q^2 + p_\perp^2}{1 - x} - \frac{p_\perp^2}{x} \right]
\]

where the gluon carries transverse momentum \( p_\perp \) and a fraction \( x \) of the initial quark momentum \( p \). For large \( p \), the emission of a gluon with moderate \( x \) and \( p_\perp \) thus requires a long formation time \( \sim 1/\Delta E \propto p \).

According to the above, typical incoming Fock states (with a small \( \Delta E \) compared to the bound state energy) should be thought of as having formed long
before the target, with a distribution among the Fock states that depends solely on the projectile wave function. The Fock states do not mix in the target region (e.g., by emitting or absorbing gluons). Hence the partons of the Fock state only scatter elastically while in the nuclear target. In particular, the parton that is to suffer the hard collision cannot lose energy by emitting gluons of moderate $p_\perp$ and $x$, if the emission time $1/\Delta E \gtrsim R_b$, the radius of the target. It may emit hard (large $p_\perp$) gluons at the hard collision vertex, but this probability is independent of the nuclear size and taken into account via higher order terms in the perturbative expansion of $\hat{\sigma}$ in Eq. (1). The quark may also emit hard gluons (such that $1/\Delta E \lesssim R_b$) somewhere else in the nucleus. This double parton scattering does depend on the nuclear size, but is suppressed by the small cross-section for two independent hard scatterings.

The same argument for limited energy loss applies to the final state. The partons that emerge from the hard collision have a typical $p_\perp \sim Q$ and hadronize into jets long after leaving the target. At the time of hadronization, the scattered parton $k$ in Eq. (1) is thus far from the other, “spectator” partons in the incoming Fock states, ensuring that the fragmentation function $D_{c/k}(z)$ is independent of the target and projectile.

A bound on the parton energy loss which follows from the above arguments (i.e., the uncertainty principle) is given in Ref. 3. It depends on the average hardness $\langle k_\perp^2 \rangle$ of the collisions that the parton experiences while traversing the nucleus, since the gluons it radiates have transverse momenta $p_\perp \leq k_\perp$. Gluons emitted in the nucleus by a parton of momentum $p$ can thus carry at most an energy fraction

$$x \lesssim \frac{k_\perp^2 R_b}{2p}$$

where $R_b$ is the target radius ($\propto A^{1/3}$ for a target of nuclear number $A$). Assuming $\langle k_\perp^2 \rangle \sim 0.1$ GeV$^2$ one finds that energy loss is negligible for high energy partons. In fact, even as its momentum $p \to \infty$ the parton suffering a hard collision can lose only a finite amount of energy in the nucleus.

Experimentally, there is convincing evidence that the energy loss of high energy quarks is insignificant (for recent reviews, see Refs. 5, 6). The evidence comes from several different processes: deep inelastic lepton scattering, lepton pair production and large $p_\perp$ jet production in $pA$ collisions. In deep inelastic lepton scattering, there are signs of energy loss when the energy $\nu$ transferred to the struck quark in large nuclei is below about 80 GeV. The $\nu$ dependence of the energy loss has been studied using models for quark hadronization inside the nucleus.

Parton energy loss has been invoked to explain the suppression of $J/\psi$ production observed at large Feynman $x$ for heavy nuclear targets. The
amount of energy loss required to explain the suppression this way is, however, considerably higher than allowed by Eq. (3). In Refs. 17, 18 and 19 a different explanation for the $J/\psi$ suppression at large $x$ is proposed, which we briefly discuss in Sect. 2.4 (see also the review of Ref. 20).

2. Spectator interactions

In the leading twist approximation the cross section depends only on single parton structure functions: The hard process is described as an incoherent sum of scatterings between one parton in the projectile and one in the target. There are kinematic situations, however, where interactions involving more than one parton in the projectile and/or in the target may be expected, and experimentally are observed, to be significant. At high energies, this is typically the case when one parton, or hadron, carries a large fraction $x$ of the energy of its parent. Next I would like to briefly discuss different aspects of these higher twist effects.

2.1 Soft Recombination with Comovers

At large scales $Q$ of the hard interaction, most partons that are involved in the hard collision have transverse momenta of $O(Q)$. Hence they do not interact significantly with the spectator partons, which have limited transverse momenta. However, a small fraction of the “hard” partons, of $O(\Lambda_{QCD}^2/Q^2)$, will be emitted with limited transverse momenta. They can then interact strongly with comoving spectators of similar velocities\textsuperscript{21}, and in particular coalesce with them to form hadrons. A specific model including such parton recombination has been studied already several years ago\textsuperscript{22}.

Parton coalescence has two principal effects:

(i) No momentum is lost in the hadronization – the hadron momentum is the sum of the momenta of the coalescing partons. Hence the $x_F$ distribution is harder than that described by the decay function $D(z)$ in Eq. (1).

(ii) There will be quantum number correlations (“leading particle effects”) between the projectile and the produced hadron.

Note that the coalescence effects are expected to be relevant only for those hadrons that are produced with a limited transverse momentum. Hence the hadron $p_\perp$ distribution should steepen with $x_F$, and only the low $p_\perp$ hadrons should show the quantum number correlations.

Evidence for comover effects have been seen, in particular, in charm hadroproduction. Since the charm quark mass is not very large, a considerable part of the charm quarks will be produced with transverse momenta small enough for their hadronization to be affected by comoving spectator quarks. Strongly enhanced
charm production at large $x_F$ has been reported by several experiments\textsuperscript{23,24,25,26,27}, but the low statistics and unconfirmed status of the observations long prevented definite conclusions. More recently, data\textsuperscript{28,29} on $\pi^- A \to D + X$ established that the $x$ distribution of the $D$ meson has the same shape as the charm quark distribution predicted by the leading twist Eq. (1). Hence the decay function $D(c \to D)$ must be assumed to be approximately $\delta(1 - z)$ in hadroproduction, significantly different from that measured\textsuperscript{30} in $e^+ e^- \to D + X$. This agrees with the comover effect\textsuperscript{31}, since the coalescence of a charm quark with a light quark of similar velocity gives a $D$ meson of momentum similar or slightly higher than the charm quark. The data\textsuperscript{28,29} is also consistent with the model of Ref. 31 for quantum number correlations between the projectile and produced $D$ meson.

It should be stressed that comover coalescence is a soft process that is not calculable in perturbative QCD. Due to the softness of the interaction, the recombination cannot significantly change the momentum of a heavy quark. Processes where, \textit{e.g.}, a fast light quark combines with a slow heavy quark, thus giving rise to a fast heavy hadron, require large momentum transfers and \textit{should} be calculated using PQCD. We shall return to this question below.

Quarkonia contain no light quarks, and thus are not formed by soft coalescence with valence quarks. On the other hand, the formation of quarkonia at moderate $x_F$ is expected to be suppressed by comovers interactions, which can break up the fragile bound state, resulting in open heavy flavor production\textsuperscript{21,31,19}. Evidence for this has been seen\textsuperscript{32} in $J/\psi$ and $\Upsilon$ production on nuclear targets, in the nuclear fragmentation region ($x_F \lesssim 0$). The suppression of quarkonium production is observed to increase with nuclear number, which is consistent with the increasing number of comoving spectators.

### 2.2 Hard Higher Twist Processes

The probability of finding two partons within a short transverse distance $1/Q$ in typical Fock states is for geometrical reasons of order $1/Q^2$. Since $1/Q$ is the coherence length of hard processes, this accounts for the size of the higher twist corrections to the single parton scattering Eq. (1). However, there are processes which only get contributions from transversally compact Fock states. In this case the corrections to the leading twist formalism can be substantial, and even dominant. A simple example is deep inelastic lepton scattering at large Bjorken $x$. A Fock state in which one quark carries $x \approx 1$ has, according to Eq. (2), a large energy and hence mixes rapidly with other Fock components. Intuitively, the partons that transferred their momentum to the leading quark must have been nearby, to accomplish the transfer in the short life-time of the $x \approx 1$ state.

More quantitatively\textsuperscript{33,20}, consider a parton that is going to transfer its longitudinal momentum, say $y p$, to the leading quark from which the lepton then scatters. This parton can have any transverse momentum up to a large value $k_\perp$ for which
the lifetime of its Fock state, \(2yp/(m^2 + k_{\perp}^2)\), becomes similar to the lifetime of the \(x \approx 1\) state, \(2p(1 - x)/(m^2 + p_{\perp}^2)\). For finite \(y\) and limited transverse momentum \(p_{\perp}\) of the \(x \approx 1\) parton, this shows that \(k_{\perp}^2 \propto 1/(1 - x)\). Hence the transverse size of the Fock states in the target hadron which can contribute to a deep inelastic scattering event with \(x \approx 1\) scales as \(r_{\perp} \propto \sqrt{1 - x}\).

Since the Fock states that participate in large \(x\) deep inelastic lepton scattering have a small transverse size, the higher twist corrections due to coherent scattering from more than one quark in the target hadron can be important. In particular, in the limit where the coherence length \(1/Q\) scales as the transverse size \(r_{\perp} \sim \sqrt{1 - x}\), the twist expansion breaks down. Higher twist contributions are not suppressed at all compared to leading twist for arbitrarily large \(Q^2\), if \(x \to 1\) as \(Q \to \infty\) such that \(\mu^2 = Q^2(1 - x)\) is held fixed.

Experimentally, the rise of the higher twist terms with increasing Bjorken \(x\) is clearly seen in deep inelastic scattering. Since the hard scattering from the transversally small target can be calculated in PQCD, the data can, in principle, be used to obtain information on multiparton correlations in the proton. The large \(x\) cross section measures the target distribution function—the probability to find all valence partons at small transverse separations, which is a function of their sharing of the total longitudinal momentum. The same distribution function also determines hard exclusive cross-sections, e.g., elastic lepton hadron scattering at large \(Q^2\).

### 2.3 Photoproduction of hadrons at large \(x_F\)

The above Fock state picture is useful also for understanding the interactions of photons. In deep inelastic scattering, the splitting of a photon of energy \(\nu\) and virtuality \(-Q^2\) into a \(q\bar{q}\) pair involves an energy difference

\[
\Delta E = -Q^2 - \frac{1}{2\nu} \frac{m_q^2 + p_{\perp}^2}{z(1 - z)}
\]

For large \(Q^2\) and \(\nu\), \(\Delta E\) is essentially independent of the fraction \(z = E_{\bar{q}}/\nu\) of the photon energy carried by the antiquark. Since a slow antiquark interacts strongly in the target, the DIS cross-section is dominated by \(z \sim \mathcal{O}(1/\nu)\). This picture of deep inelastic scattering in the target rest frame also gives an understanding of the nuclear shadowing at small \(x_{Bj} = Q^2/2M\nu\), since the transverse size of the Fock state turns out to grow like \(1/x_{Bj}\), the “Ioffe” distance from the target at which the photon creates the \(q\bar{q}\) pair. In particular, one finds that the amount of shadowing depends on the polarization of the virtual photon. This prediction has not yet been tested experimentally.
For real photoproduction, $Q^2 = 0$, an equal partition of the energy ($z \simeq \frac{1}{2}$) minimizes the energy difference $\Delta E$ in Eq. (4). These $q\bar{q}$ states are long-lived and develop into meson states before the collision – hence the photon may be represented as a mixture of vector mesons according to the Vector Dominance Model.

The pointlike photon manifests itself, however, when special requirements are imposed on the final state. Hard reactions involving the production of heavy quarks or large transverse momenta are the most well-known examples. Another, less obvious case\cite{footnote5, footnote3} is that of photoproduction of inclusive hadrons with limited $p_\perp$ but large Feynman $x_F = E_h/\nu$. As discussed above, this process also involves a short time scale. The photon can most easily create a parton with large momentum via an asymmetric $q\bar{q}$ decay, $z \propto 1 - x_F$ in Eq. (4). The quark (or antiquark) with large momentum will then penetrate the nucleus with minimal energy loss, as discussed in Section 1. This means that the large $x_F$ distribution of photoproduced hadrons will be independent of the nuclear size, contrary to the case for hadronic projectiles where an $x_F$ dependent nuclear suppression is observed\cite{footnote39}. The data on muon scattering shows\cite{footnote5} that the inclusive hadron $x_F$ distribution has the same shape for all nuclear targets, and is furthermore independent of $x_B$ and $Q^2$. In particular, the shape of the $x_F$ distribution is the same in and out of the shadowing region. This agrees with the present picture\cite{footnote3}, according to which the pointlike photon should manifest itself even at $Q^2 = 0$ when $x_F$ is large.

2.4 Heavy quark production at large $x_F$

In the standard, leading twist picture of heavy quark production represented by Eq. (1), the heavy quarks obtain their momenta from a single parton in each of the colliding hadrons. As we saw in Section 2.2, at large $x_F$ only transversally small Fock states of the projectile contribute. When the transverse size $r_\perp \propto \sqrt{1 - x}$ of the Fock state is commensurate with the Compton wavelength $1/M$ of the heavy quark, “intrinsic” diagrams, where the heavy quark pair is connected to several partons in the projectile, become important\cite{footnote33, footnote20}. Furthermore, the target parton scatters mainly from the stopped, light valence partons (which have momenta of $O(1 - x_F)$). The hardness of the interaction with the target is $\mu^2 = M^2(1 - x_F)$, which can be small even for large quark masses $M$.

So far no complete QCD calculation has been done for heavy quark production in the limit of fixed $\mu^2$. General aspects of the large $x_F J/\psi$ production\cite{footnote14, footnote15, footnote16, footnote40}, as well as model calculations\cite{footnote19}, nevertheless suggest that this limit is relevant for understanding the data. In particular, the nuclear target $A$ dependence, the decrease of the average transverse momentum of the $J/\psi$ and the Feynman scaling of the cross section all support this view\cite{footnote31, footnote19, footnote33, footnote17}, and disagree with the leading twist prediction. Furthermore, the fact that the produced $J/\psi$ is unpolarized except for $x_F \gtrsim 0.9$, where it is longitudinally polarized\cite{footnote40}, suggests a different
production mechanism at large $x_F$. An analogous change of polarization is seen in large mass lepton pair production\textsuperscript{41,42} at high $x_F$, and is understood as due to higher twist effects\textsuperscript{43}.

A particularly striking example of spectator effects is provided by a comparison of the inclusive production of $\Omega^-$ in $pN$ and $\Xi N$ collisions\textsuperscript{44}. The inclusive $pN \rightarrow \Omega^- + X$ cross-section falls steeply with $x_F$, as would be expected. On the contrary, the $\Xi N \rightarrow \Omega^- + X$ cross-section rises by more than two orders of magnitude to a maximum at $x_F \sim 0.8$. At its maximum, the $\Xi N \rightarrow \Omega + X$ cross-section is comparable to that of $pN \rightarrow \Lambda X$, and more than four orders of magnitude larger than the $pN \rightarrow \Omega$ cross-section. Clearly the “spectator” strange quarks in the $\Xi$ projectile are combining with a produced strange quark to form the $\Omega^-$. The strange quark is too light to trust the hard scattering formalism in a quantitative way. Nevertheless, the behavior of the $\Xi N \rightarrow \Omega + X$ process agrees qualitatively with what one would expect if the strange quark were truly massive. In that case, the two strange quarks in the projectile would each carry one half of the incident $\Xi$ momentum, while the three strange quarks in the $\Omega^-$ would similarly each carry $x_F/3$. Thus the incident strange quarks need to be decelerated, which requires large momentum transfers if they are considered to be heavy. A maximal $\Omega^-$ inclusive cross-section is expected when the deceleration is as small as possible, which implies $x_F \simeq .75 \ldots 1$, depending on the momentum of the $\bar{s}$ quark that is also produced in the collision.

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