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ABSTRACT

We discuss the short-range structure of hadronic and nuclear wave functions expected in QCD. In addition to the "extrinsic" contributions associated with radiation from single partons, there is an "intrinsic" hardness of the high-mass fluctuations of the wave function due to the spatial overlap of two or more partons.

We argue that intrinsically-hard partons, having large mass and/or large transverse momentum, will dominate in the region of large Feynman x_F. Their rescattering in nuclear targets is expected to be larger than for extrinsically-hard partons, leading to a suppressed production cross section for hadrons scattering on heavy nuclei. Experimental evidence for this exists for open charm, J/ψ, and Υ production at large x_F.

The effects of intrinsic hardness may be particularly striking in nuclear wave functions, where the overlap of partons belonging to different nucleons can give rise to cumulative ($x > 1$) phenomena. The data on backward cumulative particle production from nuclei supports the existence of an intrinsically-hard component in nuclear wave functions. Partons at large x_F may also be associated with the enhanced subthreshold production of particles observed in hadron-nucleus and nucleus-nucleus collisions.

We discuss the evidence for anomalies in the large angle $pp \to pp$ cross section near the charm threshold. Arguments are presented that charmonium states may bind to nuclei through the QCD Van der Waals force. This would lead to a striking signal in charm production near threshold.

INTRODUCTION

At high energies, most scattering processes only involve states that were formed long before the collision takes place. Consider the Fock expansion of a meson in QCD,

$$|h\rangle = |q\bar{q}\rangle + |q\bar{q}G\rangle + \ldots + |q\bar{q}Q\overline{Q}\rangle + \ldots \tag{1}$$

where $q(Q)$ refers to a light (heavy) quark and G to a gluon. The individual Fock components in (1) have "lifetimes" Δt (before mixing with other components)

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which can be estimated from the uncertainty relation $\Delta E \Delta t \sim 1$. At large hadron energies E the energy difference becomes small,

$$\Delta E \approx \frac{1}{2E} \left(m^2 - \sum_i \frac{m_i^2 + p_{Ti}^2}{x_i}\right) \tag{2}$$

where $x_i = (E_i + p_{Li})/(E + p_L)$ is the fractional (light-cone) momentum carried by parton i. Fock components for which $1/\Delta E$ is larger than the interaction time have thus formed before the scattering and can be regarded as independent constituents of the incoming wave function. At high energies only collisions with momentum transfers commensurate with the center of mass energy, such as deep inelastic lepton scattering ($Q^2 \sim 2m\nu$) and jet production with $p_T \sim \mathcal{O}(E_{cm})$ produce states with lifetimes as short as the scattering time.

The above arguments show that a typical scattering process is essentially determined by the mixture of incoming Fock states, *i.e.*, by the wave functions of the scattering particles. This is true even for collisions with very heavy quarks or with particles having very large p_T in the final state, provided only that the momentum transferred in the collision is small compared to E_{cm}. The cross sections for such collisions are thus determined by the probability of finding the corresponding Fock states in the beam or target particle wave functions; *cf.* Eq. (1). An example of this is provided by the Bethe-Heitler process of e^+e^- pair production in QED. A high energy photon can materialize in the Coulomb field of a nucleus into an e^+e^- pair through the exchange of a very soft photon. The creation of the massive e^+e^- pair occurs long before the collision and is associated with the wave function of the photon. The collision process itself is soft and does not significantly change the momentum distribution of the pair. Similarly, heavy quark production in hadron collisions is at high energies ($E_{cm} \gg m_Q$) governed by the hard (far off energy-shell) components of the hadronic wave functions.

The Structure of Intrinsically Hard States

The leading extrinsic contribution to heavy quarks in a hadronic wave function is one gluon splitting into a heavy quark pair, $G \to Q\overline{Q}$ (Fig. 1a). We call this contribution extrinsic since it is independent of the hadron wave function, except for its gluon content. The extrinsic heavy quarks are, in a sense, "constituents of the gluon". The extrinsic heavy quark wave function has the form

$$\Psi^{extrinsic}(q\overline{q}Q\overline{Q}) = \Gamma_G \, T_H(G \to Q\overline{Q}) \, \frac{1}{E\Delta E} \tag{3}$$

The square of the gluon amplitude Γ_G gives the ordinary gluon structure function of the hadron. The gluon splitting amplitude T_H is of order $\sqrt{\alpha_s(m_Q^2 + p_{TQ}^2)}$, and ΔE is the energy difference (2). The integral of the extrinsic probability $|\Psi^{extrinsic}|^2$ over p_{TQ}^2 for $p_{TQ} \lesssim \mathcal{O}(m_Q)$ brings a factor of m_Q^2. Hence we see that the probability of finding extrinsic heavy quarks (or large p_T) in a hadronic wave

function is actually independent of the quark mass (or p_T). This is related to the quadratic divergence of the quark loop in Fig. 1b. The production cross section of the $Q\overline{Q}$ pair is still damped by a factor $1/m_Q^2$, this being the approximate transverse area of the pair.

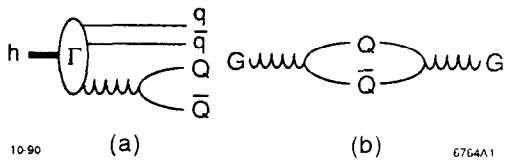

Figure 1. (a) Gluon splitting gives rise to extrinsic heavy quarks in a hadron wave function. The pointlike coupling to the gluon implies that all quark masses and all transverse momenta are generated with equal probability. (b) In the squared amplitude, this is seen as a quadratic divergence of the quark loop.

Intrinsic heavy quark Fock states[1] arise from the spatial overlap of light partons. Typical diagrams are shown in Fig. 2. The transverse distance between the participating light partons must be $\lesssim \mathcal{O}(1/m_Q)$ for them to be able to produce the heavy quarks. The wave function of the intrinsic Fock state has the general structure

$$\Psi^{intrinsic}(q\overline{q}Q\overline{Q}) = \Gamma_{ij}\; T_H(ij \to Q\overline{Q})\; \frac{1}{E\Delta E} \tag{4}$$

Here Γ_{ij} is the two-parton wave function, which has a dimension given by the inverse hadron radius. $T_H(ij \to Q\overline{Q})$ is the amplitude for two (or, more generally, several) light partons i, j to create the heavy quarks, and ΔE is the energy difference (2) between the heavy quark Fock state and the hadron. A sum over different processes, and over the momenta of the light partons, is implied in (4). In renormalizable theories such as QCD, the amplitude T_H is dimensionless. Hence, up to logarithms, the probability $|\Psi^{intrinsic}|^2$ for intrinsic heavy quarks is of $\mathcal{O}(1/m_Q^2)$ (after the p_T^2 integration). This is smaller by $1/m_Q^2$ as compared to the probability (3) for extrinsic heavy quarks,[1,2] as is true of higher twist. The relative suppression is due to the requirement that the two light partons be at a distance $\lesssim 1/m_Q$ of each other in the intrinsic contribution.

In contrast to the extrinsic contribution (3), which depends only on the inclusive single gluon distribution, an evaluation of the intrinsic Fock state (4) requires a knowledge of multiparton distributions amplitudes. In particular, we need also the distribution in transverse distance between the partons. Our relative ignorance of the multiparton amplitudes Γ_{ij} for hadrons[3] makes it difficult to reliably calculate the magnitude of the intrinsic heavy quark probability. We can, however, estimate[1] the distribution of intrinsic quarks from the size of the energy

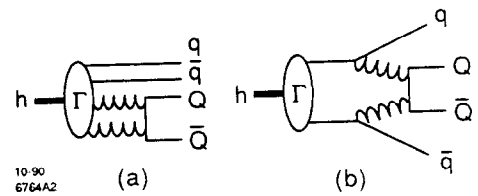

Figure 2. Intrinsic heavy quark contributions to a hadronic wave function, generated by (a) gluon fusion and (b) light quark scattering. The large mass of the produced quark implies that the participating light partons must be at a small transverse separation.

denominator ΔE, as given by (2). It is clear that those Fock states which minimize ΔE, and hence have the longest lifetimes, also have the largest probabilities. In fact, taking

$$|\Psi^{intrinsic}|^2 \sim 1/(\Delta E)^2 \ , \tag{5}$$

one finds that the maximum is reached for

$$x_i = \frac{\sqrt{m_i^2 + p_{Ti}^2}}{\sum_i \sqrt{m_i^2 + p_{Ti}^2}} \ , \tag{6}$$

implying equal (longitudinal) velocities for all partons. The rule (5) has been found to successfully describe the hadronization of heavy quarks.[5,6]

Using the probability (5), we see from (6) that partons with the largest mass or transverse momentum carry most of the longitudinal momentum. This has long been one of the hallmarks of intrinsic charm. We also note that the intrinsic heavy quark states have a larger transverse size than the extrinsic ones, although both tend to be small, of $\mathcal{O}(1/m_Q^2)$. The extrinsic heavy quarks are produced by a single (pointlike) gluon (Fig. 1), whereas the intrinsic mechanism is more peripheral (Fig. 2). This means that rescattering and absorption effects for intrinsic states produced on heavy nuclei will be relatively more significant, compared to that for extrinsic states. In addition to the heavy quarks Q, such rescattering may affect the light partons involved in the intrinsic state (*e.g.*, the quarks q in Fig. 2(b). These light quarks tend to be separated by a larger transverse distance than the heavy quarks, further enhancing the rescattering.

Consider now the formation of intrinsic heavy quark states in nuclear wave functions. At high energies, partons from different nucleons can overlap, provided only that their transverse separation is small. Thus the partons which create intrinsic heavy quarks in Fig. 2 can come from two nucleons which are separated by a longitudinal distance in the nucleus. Now it is reasonable to assume that partons belonging to different nucleons are uncorrelated, *i.e.*, that the two-parton amplitude Γ_{ij} in Eq. (4) is proportional to the product $\Gamma_i\Gamma_j$ of single parton

amplitudes. Hence the amount of intrinsic charm in nuclei may be more reliably calculated than for hadrons. The probability for intrinsic charm will increase with the nuclear path length as $A^{1/3}$. Moreover, the total longitudinal momentum of the intrinsic quark pair, being supplied by two different nucleons, can be larger than in a single hadron, and can in fact exceed the total momentum carried by one nucleon.

All that we have said above concerning heavy quark Fock states applies equally to states with light partons carrying large transverse momentum. Extrinsic and intrinsic mechanisms for generating large p_T in hadronic wave functions are shown in Fig. 3. Using Eq. (5) as a guideline for the probability of intrinsic hardness, we see in fact that the parton mass and p_T appear in an equivalent way. We again expect that the intrinsic mechanism will be dominant at large x_F, and in particular in the cumulative ($x_F > 1$) region of nuclear wave functions.

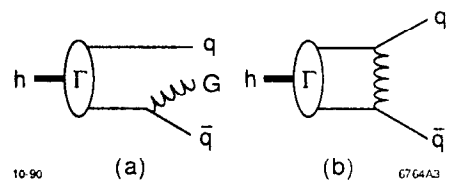

Figure 3. (a) An extrinsic contribution to large transverse momentum partons in a hadron and (b) an intrinsic contribution.

The possibility of parton fusion has been considered previously in the context of the evolution of parton distributions with momentum transfer (Q^2).[7,8] At very large Q^2 and small x, the number of gluons can become large enough to force them to overlap and coalesce. Our emphasis here is different. We are interested in rare phenomena at large x, where processes involving two or more gluons and valence quarks can give dominant effects, even though the likelihood for such fluctuations is small. The colliding partons in Figs. 1–3 are to be thought of (in a first approximation) as nearly on-shell, and having small p_T. Only the part of the processes in Fig. 3 leading to large p_T partons is to be considered as a new contribution to the wave function. In particular, the fusion of two partons into one (*e.g.*, $qG \to q$), which cannot give large p_T, is a part of the non-perturbative wave functions Γ, and hence does not contribute to intrinsic hardness.

Charm Production

The concept of intrinsic charm was originally inspired by experiments[9] showing unexpectedly abundant charm production at large $x_F = 2p_{charm}/E_{cm}$. When extrapolated to small x_F, the data suggested total charm cross sections in the millibarn range, far beyond the predictions (20 – 50 μb) of the standard QCD gluon fusion process (*cf.* Fig. 4(a)). Later data with good acceptance at low x_F showed that the total charm cross section actually is compatible with the gluon fusion process.[10] Nevertheless, more evidence was also obtained showing

that charm production at large x_F, albeit a small fraction of the total cross section, still is larger than expected.[11] The large x_F data also shows correlations (leading particle effects) with the quantum numbers of the beam hadron that are incompatible with gluon fusion.[12]

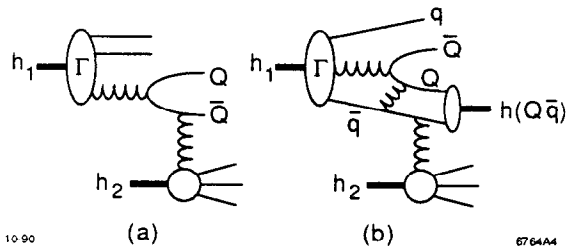

Figure 4. (a) The gluon-gluon fusion process in QCD. At high energies, the extrinsic $Q\overline{Q}$ pair preforms in the incoming wave function and is put on mass-shell by a soft gluon from the target. (b) An example of intrinsic heavy quark production. The heavy quark can get additional momentum from a light valence quark, and the produced hadrons at large x_F may get quantum numbers that are correlated to those of the valence quark (leading particle effect). The scattering can be from one of the light partons involved in the intrinsic state.

The intrinsic charm production mechanism (Fig. 4(b)) is expected to be smaller than the extrinsic one, due to the $1/m_Q^2$ suppression from the requirement of spatial overlap of initial light partons. However, at sufficiently large x_F the intrinsic mechanism will dominate, because the momentum of several incoming partons can be transferred to the heavy quarks. Our present, improved understanding of intrinsic charm, as outlined above, will allow a more quantitative theoretical discussion of these phenomena than was possible heretofore. Such an analysis will also become increasingly meaningful as the data on hadroproduced charm at large x_F improves.

Experiments on charm production from nuclear targets have shown an anomalous dependence on the nuclear number A. If the open charm (D, Λ_c) cross section is parametrized as

$$\frac{d\sigma}{dx_F} \propto A^{\alpha(x_F)} \tag{7}$$

then $\alpha(x_F \sim 0.2) \sim 0.7 \ldots 0.9$ is obtained.[10,13] For heavy nuclei ($A \approx 200$) this means a factor of $2 \ldots 3$ suppression in the cross section, compared to the leading QCD expectation ($\alpha = 1$). In this respect, the charm production data is quite different from that of massive μ-pair production, for which α is found to be very close to 1.[14]

For J/ψ production, the data on the x_F-dependence of α is particularly detailed,[15,16] showing a remarkable decrease from $\alpha = 1$ near $x_F = 0$ to $\alpha =$

0.7...0.8 at large x_F. The data at different beam energies agree, implying that Feynman scaling is valid. It is possible to show that the nuclear suppression is not due solely to the shadowing of the nuclear target structure function.[17] The effects of the target structure function can be eliminated by forming cross section ratios at a given value of the fractional momentum (x_2) of the target parton. This does not eliminate the target effects seen in the data, however, implying that the suppression does not factorize into a product of beam and target structure functions, as expected in leading twist. The target dependence thus must be due to a higher twist effect, *i.e.*, one that is of $\mathcal{O}(1/m_Q^2)$, compared to the leading (factorizable) QCD process. This is supported by preliminary data on Υ production,[16] which shows a significant but weaker nuclear suppression than for J/ψ production.

At high energies, the $c\bar{c}$ quarks do not have time to separate significantly inside the nucleus. Thus the J/ψ forms only after the charm quarks have left the nuclear environment, and the suppression cannot be related to the size of the J/ψ wave function.[18] This is also supported by preliminary data showing that the nuclear suppression for the $\Psi(2S)$ and the J/ψ is the same.[16] The $c\bar{c}$ state itself has a finite size, of $\mathcal{O}(1/m_Q^2)$, and could lose some momentum due to rescattering. Due to the rapid decrease of the cross section at large x_F, the trend of this effect is to make α decrease with x_F as observed. However, it is difficult to explain the magnitude of the x_F-dependence of α without assuming the loss of a large fraction of the momentum of the $c\bar{c}$ system.

A natural explanation of the increase of the nuclear suppression in J/ψ production with x_F is provided by the existence of two production mechanisms, the extrinsic and intrinsic ones.[19] As discussed above, intrinsic charm production is damped by a factor $1/m_c^2$, but can still dominate the small gluon fusion cross section at large x_F. Since the intrinsic heavy quark state tends to have a larger transverse size than the extrinsic one, it will suffer more rescattering in the nucleus. The x_F-dependence of α can then be understood as reflecting the increasing importance of intrinsic Fock states at large x_F.

In conclusion, the present experimental evidence for the existence of intrinsic charm is suggestive. However, the theoretical and experimental situation must improve for definite conclusions to be made. More quantitative studies of the intrinsic charm wave function, using multiparton distributions, coupled with better data on open charm at large x_F, should improve the situation in the near future.

The Intrinsic Hardness of Nuclear Wave Functions

We noted in Section 1 that intrinsic hardness should be enhanced in nuclear wave functions, due to the increased probability for spatial overlap of light partons from different nucleons. All of the data on charm production discussed above was obtained with beams of ordinary hadrons, and the experimental acceptance generally limited the observations to the forward ($x_F > 0$) hemisphere. This data thus reflects the importance of charm in the wave functions of the beam particles. An important exception to this is the EMC measurement of the charm structure

function of the Fe nucleus.[20] An enhancement over the extrinsic photon-gluon contribution was observed at large x_F, but the limited statistics prevented a firm conclusion.

Several features of scattering on nuclear targets show that the nucleus cannot always be treated as a collection of ordinary nucleons. Measurements of deep inelastic lepton scattering have revealed[21,22,23] deviations of the nuclear structure functions from those of free nucleons, both at very small and at intermediate values of x (the "EMC Effect"). There are also indications[24] that the quark distributions in nuclei extend beyond $x = 1$. Unusual states of the nucleus could be involved as well in the production of large p_T particles in hadron–nucleus collisions, where the yield is known to increase faster than the nuclear number A (The "Cronin Effect").[25,26]

The most direct evidence for an enhancement of the nuclear structure function at large x comes from the so-called "Cumulative Effect".[27] Cumulative particles are defined as hadrons produced in the fragmentation region of a nucleus which have $x_F > 1$, *i.e.*, they carry more momentum than the individual nucleons (apart from Fermi motion effects). In practice, experiments are mostly done by scattering a variety of particles (leptons, hadrons and nuclei) on stationary nuclei, and observing hadrons that are moving backward in the laboratory. A simple kinematical exercise shows that at sufficiently high beam energies, the energy E_h and longitudinal momentum p_h^L of a hadron h produced on a free stationary nucleon must satisfy

$$x \equiv \frac{E_h - p_h^L}{m_N} \leq 1 \tag{8}$$

where m_N is the nucleon mass and $p_h^L < 0$ in the backward direction. The variable x defined by (8) is the usual (light-cone) fractional momentum, which is equivalent to the Feynman momentum fraction x_F of h in the CM system. This equivalence is strictly true for infinite beam momentum; a number of alternative definitions of x have been used in order to take finite energy effects into account. The difference between the various definitions will not be important for our qualitative discussion below.

Cumulative particle production has been seen in many experiments using a variety of beam particles and energies, up to values of $x = 4$ or so. To a first approximation, Feynman scaling (*i.e.*, independence of beam energy) sets in already at quite low energies, $p_{beam} \sim 2\ GeV$ (Fig. 5(a)). The shape of the cumulative hadron distribution is insensitive to the type of beam particle used. These features suggest that the cumulative particle distribution reflects properties of the nuclear wave function.

The laboratory momenta of the cumulative particles range well beyond 1 GeV, making a description in terms of ordinary Fermi motion unlikely. If a nucleon basis is used in the wave function it would be necessary, in this energy range, to include in an essential way also N^* and Y^* excitations.[29] In fact, many

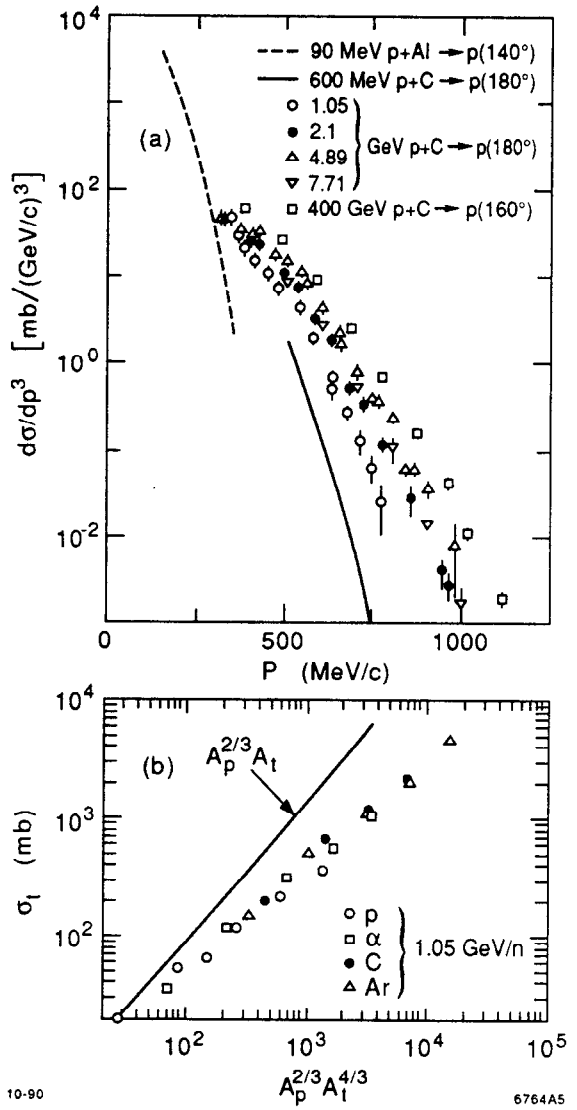

Figure 5. (a) Laboratory momentum distributions of cumulative protons produced by protons scattering on carbon and aluminum nuclei. In an analogy to the Rutherford experiment, the backscattering of 1 *GeV* protons from a beam of 2 *GeV* protons suggests encounters with small structures within the nucleus. (b) Dependence on the atomic number of the target (A_t) and projectile (A_p) for cumulative protons in the target fragmentation region. The data were fitted to a gaussian momentum distribution with a total rate parametrized by σ_t, which scales when plotted as a function of $A_p^{2/3}A_t^{4/3}$. Data and further references in Ref. 28.

arguments[27] point to the cumulative phenomena being linked to short-distance features of the nuclear wave function. The momenta of several nucleons in a nucleus have to be combined in order to produce the cumulative particles observed

at the highest values of x. This presumably requires a close spatial correlation between the nucleons. Such short-distance effects in the nuclear wave function are best described in terms of quark and gluon degrees of freedom.[30]

The dependence of the cumulative particle distribution on the atomic number of the target nucleus is at least as fast as A^1, and is compatible with $A^{4/3}$ for cumulative protons at lower energies[28,31] (Fig. 5(b)). An A-dependence this strong suggests that the production of the cumulative partons is a volume effect, with little absorption of the outgoing quanta. An $A^{4/3}$ dependence is what one would naively expect for intrinsic hardness, given that the small size of the hard cluster implies a suppression of rescattering in the nucleus, and taking into account the factor $A^{1/3}$ enhancement from the transverse overlap (of two partons) along the nuclear diameter. For nuclear projectiles, the dependence on the atomic number of the projectile is compatible[28,31] with $A_{proj}^{2/3}$. For $A_{proj} < A_{targ}$, this is also in accord with naive expectations, since the projectile presumably can put intrinsically hard clusters on their mass shell throughout a region of transverse space proportional to the area of the projectile.

Direct evidence that the cumulative phenomenon is associated with small transverse size is provided by the p_T-distribution of the produced hadrons.[32,33] The average p_T^2 grows rapidly with x, and reaches 2 GeV^2 for pions at $x = 3$ (Fig. 6(a)). This is expected for the intrinsic configurations (4), since ΔE depends on p_T^2/x (see Eq. (2)). Note that although the individual partons in an intrinsically-hard cluster (*cf.* Fig. 3) have large transverse momenta, the total transverse momentum of the cluster is small. Hence in a case such as J/ψ production, where both intrinsic quarks are incorporated in the same final hadron, much of the large p_T cancels out. On the other hand, when an intrinsic quark combines with a low p_T spectator the final hadron will carry large p_T. The experimental result that cumulative protons tend to have smaller p_T and larger cross section at a given x may be due to more intrinsic partons getting incorporated in the protons than in the pions.[30]

A remarkable feature of the cumulative x-distributions is that their shape is quite similar for all observed particles: protons, positive and negative pions and kaons. Thus, *e.g.*, the ratio between the K^- and π^- yields[34] shown in Fig. 6(b) is constant over the measured range $1.5 \leq x \leq 2.5$. This differs from the fragmentation of single nucleons,[35] for which this ratio decreases as $x \to 1$. The magnitude of the K^+ yield is[34] also much higher than would be naively expected. The heaviest nuclear targets produce roughly equal numbers of K^+ and π^+ mesons at $x \geq 1.5$.

For intrinsically-hard quarks we noted in Section 1 that the x-distribution should be similar for all quarks in a given range of p_T or quark mass, according to Eqs.(5) and (6). At the x-values considered here, the typical p_T-values are larger than, or at least comparable to, the strange quark mass (*cf.* Fig. 6(a)). Hence the π and K mesons produced by intrinsic u, d and s quarks are expected to have similar x-distributions, as observed. The K^+ mesons can get their momenta

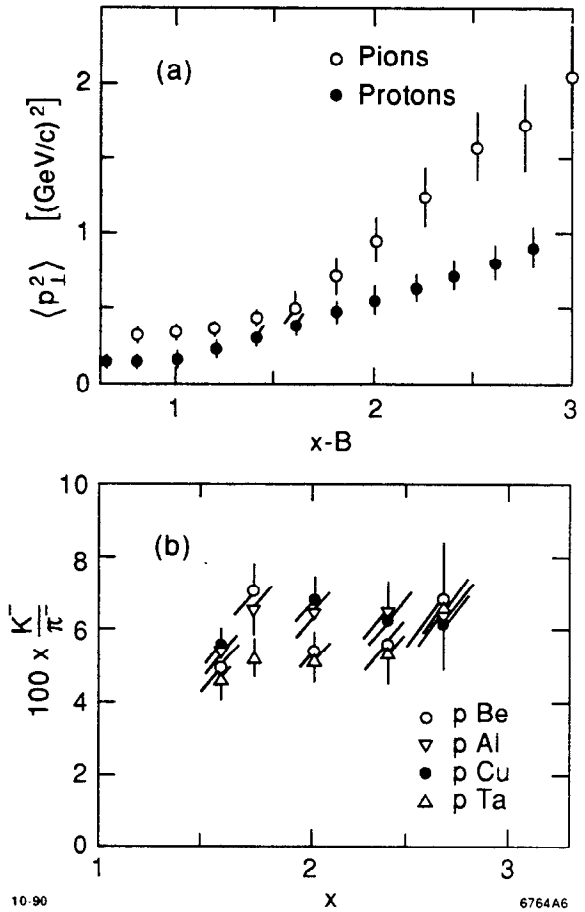

Figure 6. (a) The mean square transverse momentum of cumulative pions (o) and protons (•) produced by 10 GeV protons on Ta and Pb. The scale of the x-axis is offset by $B = 1$ for the protons ($B = 0$ for the pions). Data from Ref. 32. (b) The ratio of cumulative K^- to π^- production on several nuclei as a function of x. Data from Ref. 34.

from intrinsic u valence quarks. Since the creation of an $s\bar{s}$ pair is not suppressed at the relatively large p_T-scale involved, we can understand the equality of the K^+ and π^+ meson rates. The production of a K^- meson at large x, on the other hand, requires an energetic $\bar{u}$ or s sea quark. In this case momentum must be transferred from the valence quarks and gluons according to Fig. 2. Hence it is not surprising that the rate of K^-mesons is suppressed by about a factor 20 in the cumulative region, as seen in Fig. 6(b).

Our interpretation of the cumulative phenomena in terms of an enhancement in the nuclear structure function for $x > 1$ is compatible with some earlier suggestions.[26,27,36,37] Models of multiquark bags have been used to provide a unified explanation of the EMC, Cronin, and Cumulative Effects. An analysis of the EMC Effect in fact suggested the existence of a small admixture in nuclear wave functions of "collective" sea quarks, which are as energetic as the valence quarks.[38] The multiquark bag models do not, however, predict the probability for bag formation, nor the x-distributions of the quarks in the bag. The properties

of the intrinsically-hard component of nuclear wave functions, on the other hand, can be calculated from perturbative QCD in terms of the known quark and gluon distribution functions of nucleons. An immediate consequence is that the multiquark correlations must have a small transverse range, implying an increase of the average p_T at large x, as observed in the data (Fig. 6(a)).

Other puzzles involving fast nuclear fragments, which also may be related to intrinsic hardness, include the production of particles from nuclei below threshold for collisions on free nucleons. For example, subthreshold production of antiprotons has been observed both in $p + Cu$ and $Si + Si$ collisions.[39] While the $\bar{p}$ rate was thought to be understood for the $p + Cu$ data, based on the high cumulative momenta being interpreted as due to Fermi motion, it turned out that the corresponding calculation underestimated the rate for $Si + Si$ collisions by three orders of magnitude. In our view, the cumulative momenta should be discussed at the parton level. The rate for $\bar{p}$ production may then proceed much more favorably through, *e.g.*, the $gg \to p\bar{p}$ reaction, whose threshold is just $2m_p$ in the center-of-mass.

Effect of Heavy Quark Thresholds on Elastic pp scattering

One of the most unusual ways of identifying the effect of heavy quark thresholds is to study *elastic* scattering of hadrons at large angles. Through unitarity, even a threshold cross section of only 1 μb for the production of open charm in pp collisions will have a profound influence on the $pp \to pp$ scattering at $\sqrt{s} \sim 5\ GeV$, because of its very small cross section at 90°. The production of charm at threshold implies that there is a contribution with massive, slow-moving constituents to the pp elastic amplitude which can modify the ordinary PQCD predictions, including dimensional counting scaling laws, helicity dependence, angular dependence, and especially the "color transparency" of quasi-elastic pp scattering in a nuclear target.

It is possible to use a nucleus as a "color filter"[40,41] to separate and identify the threshold and perturbative contributions to the scattering amplitude. If the interactions of an incident hadron are controlled by gluon exchange, then the nucleus will be transparent to those fluctuations of the incident hadron wavefunction which have small transverse size. Such Fock components have a small color dipole moment and thus will interact weakly in the nucleus; conversely, Fock components with slow-moving massive quarks cannot remain compact. They will interact strongly and be absorbed during their passage through the nucleus. In fact, large momentum transfer quasi–exclusive reactions[4] are controlled in perturbative QCD by small color–singlet valence–quark Fock components of transverse size $b_\perp \sim 1/Q$; initial–state and final–state corrections to these hard reactions are suppressed. Thus, at large momentum transfer and energies, quasi–elastic exclusive reactions are predicted to occur uniformly in the nuclear volume, unaffected by initial or final state multiple–scattering or absorption of the interacting hadrons. This remarkable phenomenon is called "color transparency."[18]

Thus perturbative QCD predicts that the quasi–elastic scattering cross section will be additive in proton number in a nuclear target.[42] There are two conditions which set the kinematic scale where PQCD color transparency should

be evident. First, the hard scattering subprocess must occur at a sufficiently large momentum transfer so that only small transverse size wavefunction components $\psi(x_i, b_\perp \sim 1/Q)$ with small color dipole moments dominate the reaction. Second, the state must remain small during its transit through the nucleus. The expansion distance is controlled by the time in which the small Fock component mixes with other Fock components. By Lorentz invariance, the time scale $\tau = 2E_{\overline{p}}/\Delta\mathcal{M}^2$ grows linearly with the energy of the hadron in the nuclear rest frame, where $\Delta\mathcal{M}^2$ is the difference of invariant mass squared of the Fock components. Estimates for the expansion time are given in Refs. 41, 43, and 44.

The only existing test of color transparency is the measurement of quasi–elastic large angle *pp* scattering in nuclei at Brookhaven.[45] The transparency ratio is observed to increase as the momentum transfer increases, in agreement with the color transparency prediction. However, in contradiction to perturbative QCD expectations, the data suggests, surprisingly, that normal Glauber absorption seems to recur at the highest energies of the experiment $p_{\rm lab} \sim 12\ GeV/c$. Even more striking is that this is the same energy at which the spin correlation A_{NN} is observed to rise dramatically:[46] the cross section for protons scattering with their spins parallel and normal to the scattering plane is found to be four times as big as the cross section for anti-parallel scattering, which is again in strong contradiction to PQCD expectations.

In Ref. 47 it was noted that the breakdown of color transparency and the onset of strong spin–spin correlations can both be explained by the fact that the charm threshold occurs in *pp* collisions at $\sqrt{s} \sim 5\ GeV$ or $p_{\rm lab} \sim 12\ GeV/c$. At this energy the charm quarks are produced at rest in the center of mass, and all of the eight quarks have zero relative velocity. The eight-quark cluster thus moves through the nuclear volume with just the center-of-mass velocity. Even though the initial cluster size is small (since all valence quarks had to be at short transverse distances to exchange their momenta), the multi-quark nature and slow speed of the cluster implies that it will expand rapidly and be strongly absorbed in the nucleus. This Fock component will then not contribute to the large-angle quasi-elastic *pp* scattering in the nucleus: It will be filtered out.

The charm threshold effect will be strongly coupled to the *pp* $J = L = S = 1$ partial wave.[47] (The orbital angular momentum of the *pp* state must be odd since the charm and anti–charm quarks have opposite parity.) This partial wave predicts maximal spin correlation in A_{NN}. Hence, if this threshold contribution to the $pp \to pp$ amplitude dominates the valence quark QCD amplitude, one can understand both the large spin correlation and the breakdown of color transparency at energies close to charm threshold. Thus the nucleus acts as a filter, absorbing the threshold contribution to elastic *pp* scattering, while allowing the hard scattering perturbative QCD processes to occur additively throughout the nuclear volume.[41] One also observes a strong enhancement of A_{NN} at the threshold for strange particle production, which is again consistent with the dominance of the $J = L = S = 1$ partial wave helicity amplitude. The large size of A_{NN} observed at both the charm and strange thresholds is striking evidence of a strong effect on elastic amplitudes due to threshold production of fermion–antifermion pairs.

Nuclear Bound Quarkonium

According to the above, a slow-moving heavy quark system produced near threshold may be expected to experience strong final state interactions in the nucleus. This has interesting implications for the production of charmonium at threshold in a nuclear target. In this case it is possible that the attractive QCD van der Waals potential due to multi-gluon exchange could actually bind the η_c to light nuclei. Consider the reaction $\bar{p}\alpha \rightarrow {}^3H(c\bar{c})$ where the charmonium state is produced nearly at rest. (See Fig. 7.) At the threshold for charm production, the incident particles will be nearly stopped (in the center of mass frame) and will fuse into a compound nucleus because of the strong attractive nuclear force. The charmonium state will be attracted to the nucleus by the QCD gluonic van der Waals force. One thus expects strong final state interactions near threshold. In fact, it is argued in Ref. 48 that the $c\bar{c}$ system will bind to the ${}^3H$ nucleus. It is thus possible that a new type of exotic nuclear bound state will be formed: charmonium bound to nuclear matter. Such a state should be observable at a distinct $\bar{p}\alpha$ center of mass energy, spread by the width of the charmonium state, and it will decay to unique signatures such as $\bar{p}\alpha \rightarrow {}^3H\gamma\gamma$. The binding energy in the nucleus gives a measure of the charmonium's interactions with ordinary hadrons and nuclei; its hadronic decays will measure hadron–nucleus interactions and test color transparency starting from a unique initial state condition.

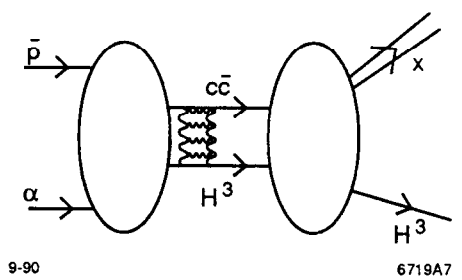

Figure 7. Formation of the $(c\bar{c}) - {}^3H$ bound state in the process $\bar{p}\alpha \rightarrow {}^3HX$.

In QCD, the nuclear forces are identified with the residual strong color interactions due to quark interchange and multiple–gluon exchange. Because of the identity of the quark constituents of nucleons, a short–range repulsive component is also present (Pauli–blocking). From this perspective, the study of heavy quarkonium interactions in nuclear matter is particularly interesting: due to the distinct flavors of the quarks involved in the quarkonium–nucleon interaction there is no quark exchange to first order in elastic processes, and thus no one–meson–exchange potential from which to build a standard nuclear potential. For the same reason, there is no Pauli–blocking and consequently no short–range nuclear repulsion. The nuclear interaction in this case is purely gluonic and thus of a different nature from the usual nuclear forces.

The production of nuclear–bound quarkonium would be the first realization of hadronic nuclei with exotic components bound by a purely gluonic potential. Furthermore, the charmonium–nucleon interaction would provide the dynamical

basis for understanding the spin–spin correlation anomaly in high energy pp elastic scattering.[47] In this case, the interaction is not strong enough to produce a bound state, but it can provide a strong enough enhancement at the heavy–quark threshold characteristic of an almost–bound system.[49]

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REFERENCES

1. S. J. Brodsky, P. Hoyer, C. Peterson, and N. Sakai, Phys. Lett. 93B, 451 (1980);
   S. J. Brodsky, C. Peterson, and N. Sakai, Phys. Rev. D23, 2745 (1981).
2. S. J. Brodsky, H. E. Haber, and J. F. Gunion, in *Anti–pp Options for the Supercollider,* Division of Particles and Fields Workshop, Chicago, IL, 1984, edited by J. E. Pilcher and A. R. White (SSC–ANL Report No. 84/01/13, Argonne, IL, 1984), p. 100;
   S. J. Brodsky, J. C. Collins, S. D. Ellis, J. F. Gunion, and A. H. Mueller, published in Snowmass Summer Study 1984, p. 227.
3. Some constraints are known for the meson distribution amplitude $\phi_M(x,Q) = \Gamma_{q\bar{q}}$ for valence $q\bar{q}$ components of the wave function. See, *e.g.*, Ref. 4.
4. S. J. Brodsky and G. P. Lepage, in *Quantum Chromodynamics*, edited by A. H. Mueller, (World Scientific, 1990.)
5. C. Peterson, D. Schlatter, I. Schmitt and P. M. Zerwas, Phys. Rev. D27, 105 (1983).
6. R. J. Cashmore, *Proc. Int. Symp. on Prod. and Decay of Heavy Flavors,* Stanford 1987 (E. Bloom and A. Fridman, Eds.), p.118.
7. L. V. Gribov, E. M. Levin and M. G. Ryskin, Nucl. Phys. B188, 555 (1981) and Phys. Rep. 100, 1 (1983);
   A. H. Mueller and J. Qiu, Nucl. Phys. B268, 427 (1986).
8. F. E. Close, J. Qiu and R. G. Roberts, Phys. Rev. D40, 2820 (1989).
9. A. Kernan and G. VanDalen, Phys. Rep. 106, 297 (1984), and references therein.
10. S. P. K. Tavernier, Rep. Prog. Phys. 50, 1439 (1987);
    U. Gasparini, *Proc. XXIV Int. Conf. on High Energy Physics*, (R. Kotthaus and J. H. Kühn, Eds., Springer 1989), p. 971.

11. S. F. Biagi, *et al.*, Z. Phys. C28, 175 (1985);
P. Chauvat, *et al.*, Phys. Lett. 199B, 304 (1987);
P. Coteus, *et al.*, Phys. Rev. Lett. 59, 1530 (1987);
C. Shipbaugh, *et al.*, Phys. Rev. Lett. 60, 2117 (1988);
M. Aguilar-Benitez, *et al.*, Z.Phys. C40, 321 (1988).

12. M. Aguilar-Benitez, *et al.*, Phys. Lett. 161B, 400 (1985) and Z. Phys. C31, 491 (1986);
S. Barlag, *et al.*, Z. Phys. C39, 451 (1988) and CERN–PPE/90–145 (1990).

13. M. MacDermott and S. Reucroft, Phys. Lett. 184B, 108 (1987);
H. Cobbaert, *et al.*, Phys. Lett. 191B, 456 (1987), *ibid.*, 206B, 546 (1988) and Z. Phys. C36, 577 (1987);
M. E. Duffy, *et al.*, Phys. Rev. Lett. 55, 1816 (1985);
M. I. Adamovich, *et al.*, CERN–EP/89–123 (1989).

14. K. J. Anderson, *et al.*, Phys. Rev. Lett. 42, 944 (1979);
A. S. Ito, *et al.*, Phys. Rev. D23, 604 (1981);
J. Badier, *et al.*, Phys. Lett. 104B, 335 (1981);
P. Bordalo, *et al.*, Phys. Lett. 193B, 368 (1987).

15. Yu. M. Antipov, *et al.*, Phys. Lett. 76B, 235 (1978);
M. J. Corden, *et al.*, Phys. Lett. 110B, 415 (1982);
J. Badier, *et al.*, Z. Phys. C20, 101 (1983);
S. Katsanevas, *et al.*, Phys. Rev. Lett. 60, 2121 (1988).

16. D. M. Alde, *et al.*, Phys. Rev. Lett. 64, 2479 (1990) and Los Alamos preprint LA–UR–90–2331 (1990);
C. S. Mishra, *et al.*, Contribution to the XXVth Rencontres de Moriond, Les Arcs (1990), Fermilab–Conf–90/100–E (May 1990).

17. P. Hoyer, M. Vänttinen and U. Sukhatme, Phys. Lett. 246B, 217 (1990).

18. A. H. Mueller, *Proc. XVII Recontre de Moriond* (1982);
S. J. Brodsky, *Proc. XIII International Symposium on Multiparticle Dynamics*, Volendam (1982);
S. J. Brodsky and A. H. Mueller, Phys. Lett. 206B, 685 (1988), and references therein.

19. S. J. Brodsky and P. Hoyer, Phys. Rev. Lett. 63, 1566 (1989).

20. J. J. Aubert, *et al.*, Nucl. Phys. B213, 31 (1983).
See also E. Hoffmann and R. Moore, Z. Phys. C20, 71 (1983).

21. J. J. Aubert, *et al.*, Phys. Lett. 123B,275 (1983);
J. Ashman *et al.*, Phys. Lett. 202B, 603 (1988);
M. Arneodo *et al.*, Phys. Lett. 211B, 493 (1988).

22. R. G. Arnold *et al.*, Phys. Rev. Lett. 52, 727 (1984).

23. G. Bari, *et al.*, Phys. Lett. 163B,282 (1985).

24. I. Savin, *Proc. XXII Conf. on High Energy Physics,* Leipzig 1984 (A. Meyer and E. Wieczorek, Eds.) Vol II, p. 251;
W. P. Schütz, *et al.*, Phys. Rev. Lett. 38, 259 (1977).

25. J. W, Cronin, *et al.*, Phys. Rev. D11, 3105 (1975);
C. Bromberg, *et al.*, Phys. Rev. Lett. 42, 1202 (1979).

26. A. V. Efremov, V. T. Kim and G. I. Lysakov, Sov. J. Nucl. Phys. 44, 151 (1986);
S. Gupta and R. M. Godbole, Phys. Rev. D33, 3453 (1986), Z. Phys. C31, 475 (1986) and Phys. Lett. 228B, 129 (1989).

27. For experimental and theoretical reviews of the cumulative effect, see:
V. S. Stavinskii, Sov. J. Part. Nucl. 10, 373 (1979);
V. B. Gavrilov and G. A. Leksin, preprint ITEP 128–89 (1989);
A. V. Efremov, Sov. J. Part. Nucl. Phys. 13, 254 (1982);
L. L. Frankfurt and M. I. Strikman, Phys. Rep. 76, 215 (1981) and Phys. Rep. 160, 235 (1988).

28. J. V. Geagea, *et al.*, Phys. Rev. Lett. 45, 1993 (1980).

29. For data on N^* production, see B. S. Yuldashev, *et al.*, WISC-EX-90-310 (1990), and references therein.

30. The fact that the number of cumulative nucleons is much larger than the number of cumulative pions (see Ref. 27) does, however, imply that the recombination of several quarks must be taken into account when using the parton basis to describe cumulative nucleons.

31. All the available data cannot, however, be fitted with a simple A^α power law dependence. See also:
Yu. D. Bayukov, *et al.*, Phys. Rev. C20, 764 (1979) and Sov. J. Nucl. Phys. 42, 238 (1985);
S. Frankel, *et al.*, Phys. Rev. C20, 2257 (1979);
N. A. Nikiforov, *et al.*, Phys. Rev. C22, 700 (1980);
M. Kh. Anikina, *et al.*, Sov. J. Nucl. Phys. 40, 311 (1984);
G. R. Gulkanyan, *et al.*, Sov. J. Nucl. Phys. 50, 259 (1989).

32. S. V. Boyarinov, *et al.*, Sov. J. Nucl. Phys. 46, 871 (1987).

33. A. I. Anoshin, *et al.*, Sov. J. Nucl. Phys. 36, 400 (1982);
A. M. Baldin, *et al.*, Sov. J. Nucl. Phys. 39, 766 (1984);

34. S. V. Boyarinov, *et al.*, Sov. J. Nucl. Phys. 50, 996 (1989).

35. P. Capiluppi, *et al.*, Nucl. Phys. B79, 189 (1974).

36. A. M. Baldin, Nucl. Phys. A434, 695 (1985).

37. A. V. Efremov, Sov. J. Nucl. Phys. 24, 633 (1976);
G. Berlad, A. Dar and G. Eilam, Phys. Lett. 93B, 86 (1980);
C. E. Carlson, K. E. Lassila and U. P. Sukhatme, preprint WM-90-115 (1990).

38. A. V. Efremov, Phys. Lett. 174B, 219 (1986);
A. V. Efremov, A. B. Kaidalov, V. T. Kim, G. I. Lykasov and N. V Slavin, Sov. J. Nucl. Phys. 47, 868 (1988).

39. J. B. Carroll, *et al.*, Phys. Rev. Lett. 62, 1829 (1989);
A. Shor, *et al.*, Phys. Rev. Lett. 63, 2192 (1989), and references therein.

40. G. Bertsch, S. J. Brodsky, A. S. Goldhaber, and J. Gunion, Phys. Rev. Lett. 47, 297 (1981).
41. J. P. Ralston and B. Pire, Phys. Rev. Lett. 61, 1823 (1988), University of Kansas preprint 90–0548 (1990).
42. By definition, quasi–elastic processes are nearly coplanar, integrated over the Fermi motion of the protons in the nucleus. Such processes are nearly exclusive in the sense that no extra hadrons are allowed in the final state.
43. B. K. Jennings and G. A. Miller, Phys. Lett. B236, 209 (1990). and University of Washington preprint 40427-20-N90 (1990).
44. G. R. Farrar, H. Liu, L. L. Frankfurt, M. I. Strikman, Phys. Rev. Lett. 61, 686 (1988).
45. A. S. Carroll, *et al.*, Phys. Rev. Lett. 61, 1698 (1988).
46. G. R. Court, *et al.*, Phys. Rev. Lett. 57, 507 (1986).
47. S. J. Brodsky and G. de Teramond, Phys. Rev. Lett. 60, 1924 (1988).
48. S. J. Brodsky, G. de Teramond, and I. Schmidt, Phys. Rev. Lett. 64, 1011 (1990).
49. The signal for the production of almost–bound nucleon (or nuclear) charmonium systems near threshold is the isotropic production of the recoil nucleon (or nucleus) at large invariant mass $M_X \simeq M_{\eta_c}$.