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Color Transparency Phenomenon and Nuclear Physics

Measuring color transparency effects provides a promising new method to investigate nucleon and nuclear structure in the domain of non-perturbative QCD, construct a theory of superdense nuclear matter, and disentangle the physics of heavy ion collisions. The simplest versions of popular hadronic models show that interesting pheonomena may exist for momentum transfers as low as about 1 GeV/c^2.

Key Words: *color transparency, QCD, final state interactions, point-like configurations*

1. INTRODUCTION

Final state interactions, FSI, are the bane of a physicist's existence. This Comment is concerned with the possibility that FSI do not occur in certain high momentum transfer reactions involving nuclear targets. This absence is caused by the cancellation of color fields produced by a system of closely separated quarks and gluons, and is termed color transparency. We will explain how the observation of this phenomenon would become a new testing ground of models of nucleons, nuclei, and QCD. The physics we consider is accessible to lepton and hadron beams of energy from 4 to 40 GeV/c. Related physics issues at larger energies are discussed in Ref. 1.

The absence of FSI may seem very surprising to some. We therefore begin with an old example. Consider the decays of ul-

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trafast pions $\pi^0 \rightarrow e^+e^-\gamma$ in emulsion and study the dependence of ionization on the distance traveled by the e^+e^- pair. Following a suggestion by King, Perkins[2] found that for small distances there are very few ionizing interactions (Fig. 1). Perkins explained: initially the e^+ and e^- are produced at the same point, so the pair acts as a dipole with small radius. The electromagnetic interactions are cancelled and there are very few ionizing interactions. The initial small ($r \sim r_\pi$) pair is not a positron eigenstate, so it separates as it moves. Thus at larger separations, interactions do occur. The data sample used by Perkins involved only 7 events, but further cosmic ray measurements confirmed this charge screening effect.[3,4]

Such charge screening effects may occur in QCD as effects of color transparency. The existence of color transparency depends on (1) the formation of a small-sized wave packet in a high momentum transfer reaction. (2) The interaction between such an object and nucleons is suppressed (color neutrality or screening) in a manner similar to the example of the previous paragraph. (3) If the wave packet escapes the nucleus while still small, no or reduced FSI occur. The remainder of this section defines the term "color transparency" and explains the necessary three requirements. Examples of color screening phenomena are discussed in Section 2. Section 3 is concerned with how the wave packet escapes from the nucleus. The ability to form a small wave packet is discussed in the context of constituent quark, bag, cloudy bag and

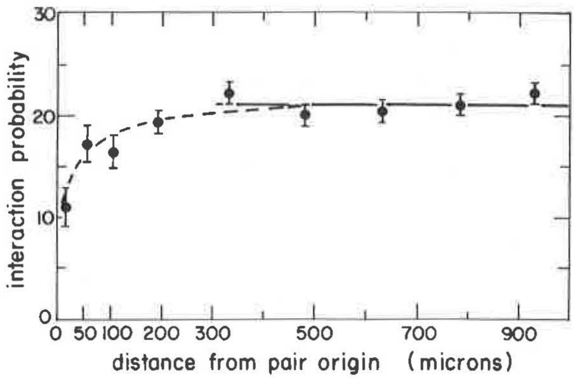

FIGURE 1 Interaction (ionization) probability vs. distance from the origin of the e^+e^- pair.

Skyrmion models and lattice theory calculations in Section 4. The effects of small-sized wave packets on the nuclear ground state and heavy ion collisions are also discussed in Section 4. The scant relevant data are reviewed in Section 5. The final section contains a summary and perspectives for the future.

1.1. Defining Color Transparency

First we discuss the notion of transparency in nuclear physics. In traditional nuclear theory, strongly interacting particles experience absorbtion. Imagine shooting a beam of protons at a nucleus, and measuring the number that emerge with momentum close to the initial value. The number coming out is smaller than the number coming in by a factor $\sim e^{-L\rho\sigma}$. Here L is the distance traversed by a particle moving straight through the nucleus, ρ is the nuclear density, and the proton–nucleon total cross section is σ. The exponential decrease with L and the corresponding forward peaked Frauenhoffer diffraction pattern nature of the angular distribution for elastic scattering are the signatures of opacity or black disk physics.

Now consider another situation. Suppose a proton impinges on a nucleus, and knocks out a proton initially at rest, in a high momentum transfer reaction. The idea of color transparency is that in this case there is no exponential loss of flux—the black disk is not present and *transparency* exists!

This is just one example, the $(p, 2p)$ reaction, of color transparency. This phenomenon could occur in a number of high momentum transfer reactions. Others could be the wide angle high energy $(e, e'p)$ or $(\pi, \pi p)$ reactions.

To be specific, we consider processes for which the fundamental reaction is elastic, or at least a two-body reaction. For example, in electron scattering the proton absorption of a high energy virtual photon leads to a high energy proton, or a Δ, or an N^*. Then a high energy baryon emerges from the nucleus, with the residual nucleus often left in an excited state. One must know the excitation energy to ensure that no extra pions are created in the hadronic process. Reactions in which the kinematic resolution is good enough to assure that the hadronic reaction is two body have been termed "semi-exclusive".[5] Here we define color transparency CT as the suppression of absorption in hadron–nucleus "semi-exclusive" interactions.

1.2. Origins of Color Transparency

We discuss the three requirements for color transparency.

1. Small objects are produced at high Q^2 in two-body wide angle reactions. High momentum transfer is associated (via Fourier transformations) with small lengths. Here we go a step further. We assert that if a proton absorbs a photon of four-momentum transfer squared $-Q^2$, the ejected wave packet (ejectile) as well as dominant configurations in the initial protons behave as a small object of transverse size $1/Q^a$ where $1 \geq a > 0$.

To see how the small size arises, consider first an example based on quark counting rules. Suppose a quark in a proton absorbs a photon. The struck quark is off energy shell by $\Delta E \sim |\mathbf{q}| = Q$, so it has a lifetime τ determined by the uncertainty principle $\tau = 1/Q$. The virtual quark must decay by emitting a particle (gluon). We want to measure a proton (or low-lying nucleon isobar) in the final state. The final state can be a proton (isobar) only if nearby quarks absorb the decay products; the other quarks must be a distance $r \simeq ct \sim 1/Q$ away from the off-shell quark.

This argument has the flavor of perturbative quantum chromodynamics pQCD, but ignores the infinite number of bare quarks and gluons that occupy each hadron. A more general argument is that at high energies, in a wide angle high energy two-body collision, the color current abruptly changes its direction. Such a process is ordinarily accompanied by radiation of gluons and the formation of multi-particle final states. The only possibility to avoid this radiation (and have a two-body final state) is if the color charge is localized in a small volume before and after the collision. (An accelerated point-like neutral object does not radiate.) Thus the interacting hadronic configurations need not be those of a minimal number of partons. Furthermore the Feynman mechanism[6] (in which a single quark carrying the whole hadronic momentum is turned around by the hard photon) becomes inoperative at large Q^2, since too many gluons would be emitted. This gluon radiation physics is strictly valid for sufficiently large momentum transfer and the size $r \sim 1/Q^a$ with $(0 < a < 1)$ instead of the $1/Q$ of quark counting rules.

An interesting question concerns the value of Q^2 required for the gluon radiation effects to be strong enough to suppress the Feynman mechanism. This is important, since in the Feynman mechanism the hadron size arises from the transverse distribution of slow or wee partons. Since the wee partons are not influenced by the hard collision, a small sized ejectile would not be formed and color transparency would not occur. The current level of understanding of QCD does not allow us to know the minimum momentum transfer required. If radiation of gluons with transverse momenta as small as 0.3 GeV/c is essential in the hard processes, then values of Q^2 as small as several GeV^2 would be sufficient. This scenario is popular in the theoretical description of hard processes, see, e.g., Ref. 7. However, in the dispersion sum rule approach, in which it is argued that perturbative QCD effects are a small correction, much larger values of Q^2 are required. For example, Bakulov and Radushkin[8] find the Feynman mechanism to dominate the pion form factor for values of Q^2 up to 10 GeV^2. Thus if the dispersion sum rule approach to the pion (nucleon) form factor is correct, color transparency may not be observed in experimentally feasible nuclear quasielastic processes.

The small object created in a high Q^2 reaction is a coherent superposition of physical states. This is a wave packet. Some authors use the expression "little proton" to describe this wave packet. This is not correct: the only proton is the physical proton with its well measured size and trivial eigenstate time dependence. We use the term point-like transverse sized configurations, PLC here, though small size configuration, SSC, may be more appropriate. Other authors[9,10] term the small object an "ejectile".

2. *Small objects have small cross sections.* The color field of a small color neutral (singlet) object is suppressed because fields of individual closely separated quarks and gluons cancel each other. Thus the interactions between the small color singlet wave packet and nucleons are smaller than those between ordinary hadrons and nucleons. This is a consequence of color neutrality which is the QCD analog of the QED concept of charge neutrality. Another popular name for color neutrality is color screening. We treat these

two phrases as equivalent. Color screening also predicts that hadronic cross sections depend on hadronic sizes, with hadron (h)–nucleon (N) cross sections varying as $\sigma_{hN} \sim r_h^2 r_N^2$. This relation is called *geometric scaling* which seems to be applicable in a larger kinematical region than pQCD. As discussed in Section 2, there is a great variety of evidence supporting the idea that small objects have small cross sections.

3. The PLC must escape the nucleus before expanding. The PLC is not a stationary state, so it undergoes time evolution. Its size increases, since (by definition) the PLC starts as small. Consider a small object in its rest frame. It expands with a characteristic time defined as τ_0. A reasonable guess is that $\tau_0 \approx 1$ fm. Now suppose the object moves with high energy E in the lab. Time dilation increases this time by a factor of E/m, so the relevant expansion time τ is given by $\tau = E/m\tau_0$. For sufficiently large energies, τ is long enough so that the object can leave the nucleus while small enough to avoid final state interactions. Then CT occurs. References 1, 5, 9, and 10 contain arguments that the mass m used to compute the time dilation effect is about equal to the average of the masses of the nucleon and its first excitation. For presently available energies, $\tau \lesssim 5\tau_0 \approx 5$ fm is small enough so that the PLC expands significantly as it moves through the nucleus. Thus the final state interaction is suppressed but does not disappear.

CT means that only the small size configurations get through the nucleus. These PLC have overlaps with baryon resonances. Thus a large cross section for baryon resonance production is expected, even for those resonances which ordinarily interact strongly with the nucleons.[11,12] This is the analogue of diffraction of light from a black disk.

The idea that a hadron in a PLC may interact with a small cross section was discussed a long time ago[67] as an explanation of the smallness of the $J\Psi$-nucleon cross section and was illustrated by an analysis of pair-production by high energy photons in the nuclear Coulomb field.[68]

The requirements for color transparency to occur in nuclear physics were set down by Mueller[13] and Brodsky[14] on the basis of perturbative QCD. The importance of the expansion of the PLC with time was found by Farrar *et al.*[5] and by Jennings and Miller.[9,10,15]

(See also Ref. 16.) The idea that color neutrality (screening) causes the cross section for the interaction of color singlet configurations with hadrons to be proportional to the square of the radius of the region occupied by color was put forward by Low,[17] Nussinov[18] and Gunion and Soper.[19] These authors assumed the high energy hadron–hadron interaction to be dominated by the exchange of two gluons. The applications of this model to soft high energy hadron collisions and references are discussed in Ref. 20. The limitations of this two gluon exchange model are discussed in Section 2. It is interesting to note that color neutrality (screening) is closely connected with the observation of Bjorken scaling at small x; see Refs. 1, 21, 22 and Section 2 below.

A nonperturbative mechanism for suppression of the interaction of PLC with hadrons has been suggested by Frankfurt and Strikman[23] on the basis of quark models: the emission of mesons from a baryon in a PLC is suppressed because the meson–baryon wave function overlap is small if the baryon is very small. Indeed the one-meson exchange contribution to the ejectile–nucleon scattering amplitude is expected to have a dependence on size similar to that obtained from the two-gluon exchange model.

We believe it desirable to use the separate terms color screening and color transparency. The term color transparency describes the vanishing of the hadron–nucleus interactions in semi-exclusive high Q^2 reactions. Color screening leads to geometric scaling and Bjorken scaling, so it seems well verified. Unlike color screening, the existence of color transparency at feasible energies is still an open question.

2. COLOR SCREENING (NEUTRALITY) PHENOMENA

Quantum chromodynamics (QCD) is a non-Abelian gauge theory of quarks and gluons. The most prominent property of such a theory is that of asymptotic freedom. We are concerned here with the other distinctive feature of QCD-color screening.

We start with the evidence concerning geometric scaling. The cross section for the interaction of an energetic colorless wave packet of small transverse size PLC characterized by a length b with a target should be small in QCD:

$$\sigma^{PLC} \sim \pi b^2. \qquad (1)$$

At sufficiently small b^2 the running coupling constant $\alpha_s(b^2)$ is small enough to apply methods of perturbative QCD. Then two gluon exchange effects are dominant, except for the processes occurring at extremely high energies. Equation (1) is a consequence of color screening and gauge invariance in the two-gluon exchange model,[17–19] if $\alpha_s(b^2) << 1$ and the beam energy is large enough so that the configuration is unchanged (frozen) during the collision. Logarithmic correction ($\ln(b^2)$) terms are ignored in Eq. (1).

The two gluon exchange model of soft hadron processes can be used as an illustrative model only since the coupling constant is not truly small. Furthermore, the calculated radiative corrections to this model lead to an increase of the cross section as $\approx s^{1/2}$,[24] which is obviously at variance with the experimental data. The current experimental data (e.g., Ref. 25) on diffractive hadron production corresponds to behavior in between $(1 - x)^5$ and $x(1 - x)$.

It is worthwhile noting that Eq. (1) is not the same as the standard one ($2\pi(r_1 + r_2)^2$) for collisions of two black objects of radii r_1 and r_2. In the standard formula, the cross section does not vanish as r_1 tends towards 0. But color screening effects can completely eliminate a cross section.

Color screening has been investigated many times. The most important manifestation occurs in the existence of precocious Bjorken scaling in the deep inelastic lepton–proton scattering (DIS). At low x, the interaction proceeds by the conversion of the highly virtual γ into a $q\bar{q}$ well before the photon hits the target. The lifetime of the $q\bar{q}$ fluctuation is $1/2m_N x$, which is large for small values of x. Thus the small x DIS process is the QCD analogue of the Perkins experiment.[2] For scaling to be obtained, the large $q\bar{q}$–proton interaction must be restricted to configurations in which the pair has a small transverse momentum[22] and hence a significant transverse separation. Another observation is that including color screening effects prevents the cross section for real photons from increasing rapidly with energy at FNAL energies.[1]

Still another method[1,21] is to study jet production in deep inelastic lepton scattering off nuclei. This is because nuclear shadowing depends on the intrinsic transverse momentum of produced jets. The transverse momentum κ_t varies as $\kappa_t^2 \sim 1/b^2$, so that jets characterized by large values of κ_t should undergo less nuclear shadowing than jets of smaller κ_t. One more consequence of color screening occurs in the nuclear photoproduction of charmed particles. The essential distance between the c and $\bar{c}$ quarks is $b^2 \sim 1/m_c^2$ which is much smaller than for the lightest quarks. Thus, color screening should be more essential in this case[1,21] and there should be less nuclear shadowing than for the production of particles made of lighter quarks.

The analysis[1,21] shows that experimental data on the phenomena mentioned in the preceding paragraph are all consistent with the notion that Eq. (1) may be applicable for rather large values of b^2:

$$b^2 \lesssim 0.25 \text{ fm}^2. \tag{2}$$

Moreover the observed pattern of total cross sections of hadron–nucleon scattering

$$\frac{\sigma_{\pi N}}{\langle r_\pi^2 \rangle} \simeq \frac{\sigma_{pp}}{\langle r_p^2 \rangle} \simeq \frac{\sigma_{KN}}{\langle r_K^2 \rangle} \tag{3}$$

is entirely consistent with Eq. (1).[26]

The appearance of Eq. (1) and its gluonic origin might seem rather strange to most nuclear physicists who take interactions to originate from mesonic exchanges. It turns out that an equation like Eq. (1) holds for the emission of a meson (M) from a PLC. The radiation of a meson from a nucleonic (N) PLC is suppressed by the factor:

$$R = \frac{\Gamma_{plc,plc,M}/P_{plc}}{\Gamma_{N,N,M}} \approx (r_{plc}/r_M)^2. \tag{4}$$

Here Γ is the appropriate vertex; P_{plc} is the probability that the point-like small size configuration exists in the nucleon, and r_M is the average radius of the meson. Studies[1,21,23,27] of the coupling of a three-quark PLC to a composite meson show that Eq. (4) is reasonable.

High energy diffractive processes also provide evidence for geometric scaling. At Fermilab and ISR energies, wave packets have no significant expansion during the reaction process, so one may concentrate on the b^2-dependence of the wave packet–nucleon scattering amplitude. This is mainly imaginary, so the optical theorem allows us to speak of the wave packet–nucleon cross section. The relevant quark distribution can be characterized as $\Psi^2(b^2)$, the square of the hadron wave function integrated over the longitudinal momentum, with b the transverse size operator. The cross section depends upon b^2 according to, e.g., Eq. (1). Thus the hadron–nucleon hN scattering cross section is given by

$$\sigma_{hN} = \int d^2b\Psi^2(b^2)\sigma(b^2). \tag{5}$$

It is useful to rewrite Eq. (5) in terms of a probability, $P(\sigma)$, that the system has a cross section, σ. Then

$$\sigma_{hN} = \int d\sigma P(\sigma)\sigma. \tag{6}$$

The quantity $P(\sigma)$ combines information regarding the basic wave function (space size, number of wee partons, etc.) and the b^2 dependence of the scattering operator in a manner that can be measured in diffractive processes. The central issue is the width of $P(\sigma)$. For example, it would be difficult to see how Eq. (1) could be valid if $P(\sigma)$ were a delta function.

An analysis[28] of FNAL and ISR data shows that $\int d\sigma P(\sigma)\sigma^2 = 1.2 - 1.3\sigma_{hN}^2$, so that $P(\sigma)$ is broad. Practically the same value for the second moment follows from the analysis of the inelastic correction to the total p-deuteron cross section.[29] This connection was first noticed in Ref. 30. The broad nature of $P(\sigma)$ supports the notion of color screening, Eq. (1) and the existence of PLC.

3. TIME SCALE FOR ESCAPE

Suppose a PLC is produced in the interior of the nucleus. Let τ_0 be the rest-frame time for a quantum fluctuation to take the PLC into a normal sized configuration. A moving PLC expands in a time $\tau = \tau_0 E/m$ where the factor E/m accounts for time dilation effects. At sufficiently large energies E, the length ($= c\tau$) required for expansion can be larger than the nuclear diameter, and the PLC leaves the nucleus without much expansion. In that high E situation the PLC can be regarded as "frozen" during the escape process.

At lower energies, the PLC does expand as it moves through the nucleus. The need to incorporate the effects of this expansion in calculations of nuclear cross sections was recognized by Farrar *et al.*[5] Those authors model the PLC–nucleon cross section as a function of the distance travelled Z from the point of hard interaction (cf. Fig. 1):

$$\sigma_N^{\mathrm{PLC}}(Z) = \left(\sigma_{\mathrm{hard}} + \frac{Z}{2p/\mu^2}\left[\sigma - \sigma_{\mathrm{hard}}\right]\right)\theta(2p/\mu^2 - Z) + \sigma\theta\left(Z - \frac{2p}{\mu^2}\right). \quad (7)$$

Here $\mu^2 \sim 0.5–0.6\ \mathrm{GeV}^2$ as determined from multiperipheral processes at high energies.[21]

The linear dependence on Z can be seen by considering DIS at small x for which the process $\gamma \rightarrow q\bar{q}$ dominates. The $q\bar{q}$–nucleon cross section must vary as $\sigma \propto 1/Q^2$, if Bjorken scaling is to be obtained. But the $q\bar{q}$ pair has a lifetime $1/2m_N x = \nu/Q^2$. Equating the distance Z propagated with the lifetime yields $Z \propto \nu/Q^2 = p\sigma$, which is the essential feature of Eq. (7). Thus the underlying physics behind Eq. (7) is that at the interaction point the square of the effective mass of the wave packet is large, $\sim Q^2$, so the expansion rate is governed by the Lorentz dilated lifetime $\sim p/Q^2$ which is much smaller than p/m_h^2, while the interaction cross section for this wave packet $\sim 1/Q^2$.[21] In pQCD this reasoning can be checked by analyzing the relevant Feynman diagrams (see the discussion in Ref. 1).

The time development of the PLC can also be computed using hadronic basis states.[9,10,12,15,16] Start with the PLC as formed by a hard interaction T_H acting on a nucleon:

$$|\mathrm{PLC}\rangle = T_H|N\rangle. \tag{8}$$

The projection of $|\mathrm{PLC}\rangle$ on a nucleon $|N\rangle$ is the form factor $F(Q^2)$. The PLC can be treated as a coherent linear superposition of hadronic states that propagate through the nucleus. The interaction with the nucleus U depends on the transverse separation b^2 between any pair of quarks in the $|\mathrm{PLC}\rangle$. In the optical approximation

$$U = -i\sigma(b^2)\rho(R). \tag{9}$$

Only the imaginary part of U is kept and the optical theorem is used. The density of target nucleus is $\rho(R)$. The amplitude, $\mathcal{M}$, for the reaction is given as

$$\mathcal{M} = \langle N|(1 + UG\cdots)|\mathrm{PLC}\rangle. \tag{10}$$

At this stage we express the $|\mathrm{PLC}\rangle$ in terms of baryonic states, m. The Green's function G is replaced by a sum of baryonic propagators G_m. Then to first order in U,

$$\mathcal{M} \approx F_N(Q^2) + \sum_m \int e^{-ipZ}\langle N|U|m\rangle e^{ip_mZ}\langle m|\mathrm{PLC}\rangle\cdots \tag{11}$$

where p is the momentum of the outgoing nucleon and $p_m \approx p + (M^2 - M_m^2)/2p$. The exponential e^{ip_mZ} arises from the eikonal propagator G_m. For later purposes it is convenient to define operators $\hat{U}$ and $\hat{p}$ ($\hat{p}^2|m\rangle \equiv (p^2 + M_m^2)|m\rangle$) so that

$$\hat{U} \equiv e^{-i\hat{p}Z}Ue^{i\hat{p}Z}. \tag{12}$$

Consider the work of Jennings and Miller as a concrete example. They use Eq. (1) to describe $\sigma(b^2)$ and a two-dimensional harmonic oscillator for the baryonic states m. Then b^2 connects the nucleon to the nucleon and only one resonant state (of mass M_1). The first order result[31] is that the σ of a standard calculation is replaced by σ_{eff}:

$$\sigma_{eff}(Z) = \sigma(1 - e^{-i\Delta p Z}) \tag{13}$$

where $\Delta p = (M_1^2 - M^2)/2p$ and Z is the propagation length. Using the frozen approximation of $p \to \infty$ leads to $\sigma_{eff} = 0$, and color transparency occurs.

The quantity σ_{eff} is complex because it is derived from a calculation of a scattering amplitude. The contribution to $\mathcal{M}$ proportional to Im σ_{eff} is not coherent with the Born term and therefore has little numerical effect for values of M_1 ranging from 1.3 to 2 GeV. The real part of $\sigma_{eff} \sim Z^2$ for small values of Z instead of the linear dependence of Eq. (7). Such a dependence could arise from a complex expansion coefficient present in a more general treatment of the |PLC⟩. The $(e, e'p)$ and $(p, 2p)$ cross sections predicted by Ref. 5 are only somewhat different from those of Refs. 10 and 20. The reason is shown in Fig. 2 which presents a comparison of σ_{PLC} with the Re σ_{eff}. (In each case the PLC is taken as having no initial size.) For the displayed important values

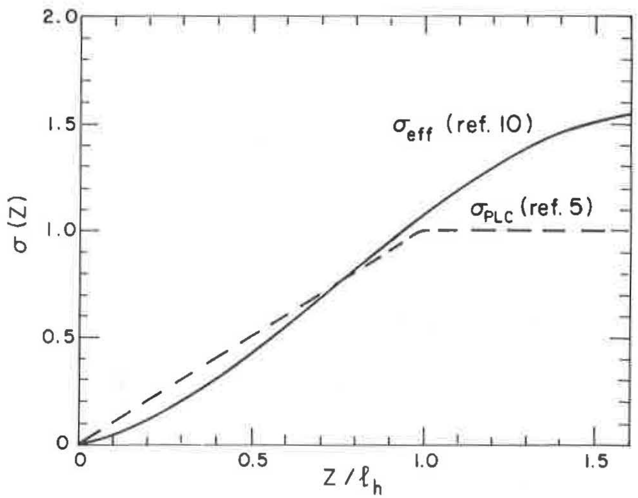

FIGURE 2 Comparison of σ_{eff}^{solid} of Ref. 9 with σ_{PLC} of Refs. 1 and 19. The distance that the PLC travels is Z and $l_h \equiv 2p/\mu^2$.

of Z, the two terms are about the same. For larger values, the PLC is no longer small and the nuclear absorption prevents much further propagation.

The above discussion involves a first order treatment of U. It is thus worthwhile to consider an exactly solvable model in which the PLC is described as a superposition:

$$|\mathrm{PLC}\rangle = \alpha|N\rangle + \beta|N^*\rangle. \tag{14}$$

In general α and β are complex numbers representing elastic and inelastic form factors. (In Ref. 10, $\alpha/\beta = 1$.) Instead of using a specific function $\sigma(b^2)$ to construct the interaction $\hat{U}(Z)$ as in Eqs. (9) and (12) we demand that

$$\hat{U}(Z = 0)|\mathrm{PLC}\rangle = 0. \tag{15}$$

In the two-state basis this means that

$$\hat{U}(Z) = -i\sigma\rho\begin{pmatrix} 1 & -\frac{\alpha}{\beta}\, e^{i\Delta pZ} \\ -\frac{\alpha}{\beta}\, e^{-i\Delta pZ} & \alpha^2/\beta^2 \end{pmatrix}. \tag{16}$$

Here color screening Eq. (15) is incorporated via modeling a set of matrix elements. Note that the ratios of the soft amplitudes for N^* to N production in nucleon–nucleon scattering are equal to ratios of hard form factors. This is an example of a deep relationship between soft and hard processes that should exist in QCD. The necessary many states would complicate a more realistic treatment. It may be easier to examine the $c\bar{c}$ system $(J/\psi, \psi')$[1] to discover these soft–hard relationships.

The interactions of $\hat{U}$ of Eq. (16) vanish at $Z = 0$, increase with Z to a maximum at $\Delta pZ = \pi$ and vanish again when $\Delta pZ = 2\pi$. Thus the size of the PLC changes as it moves through the nucleus. The transition operator T is obtained by solving the equation

$$2i\,\frac{\partial T(Z)}{\partial Z} = \hat{U}T(Z). \tag{17}$$

The model is rendered exactly soluble by assuming that the PLC is produced at $Z = 0$ and propagates through a nuclear slab of length L. Solving for the operator

$$\begin{pmatrix} e^{-i\Delta p Z} & 0 \\ 0 & 1 \end{pmatrix} T$$

simplifies the remaining algebra. The amplitude is

$$\mathcal{M} = \langle N|T(L)|\text{PLC}\rangle, \tag{18}$$

with the probability of a nucleon to traverse a distance L without energy loss given by $|\mathcal{M}|^2$. Similarly, the N^* production $\mathcal{M}_{N^*}$ amplitude is

$$\mathcal{M}_{N^*} = \langle N^*|T(L)|\text{PLC}\rangle. \tag{19}$$

The toy model results depend on the parameters M_1 and α/β. An example is shown in Fig. 3 with $M_1 = 2.4$ GeV, $\alpha/\beta = 0.8$ and $\sigma\rho = 25$ mb (1/6 fm^{-3}). The displayed oscillations occur for many values of α/β. Color transparency at high energies generally occurs.

4. PLC IN HADRONS

Hadronic wavefunctions consist of many configurations, each with its own coordinate space distribution. It is reasonable to expect that high momentum transfer reactions are influenced by configurations of much smaller size than average. A natural candidate is a configuration with the minimal number of bare quarks. The masses of bare u, d, s quarks are small, so to avoid a ghost pole due to transition of a nucleon into a system of three bare quarks, the transverse momenta of quarks should be large. Thus the minimal Fock space configuration has a small size.[1,21] This idea is supported by the recent lattice study[32] of the pion wave function. Figure 4 shows that the size of the $q\bar{q}$ component is about half of the pion size.

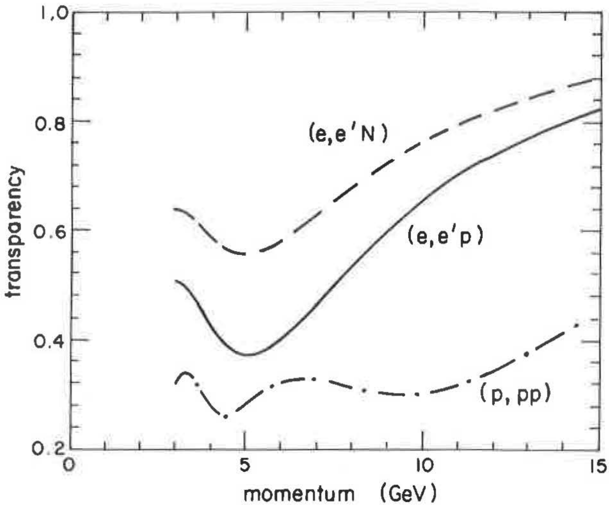

FIGURE 3 Exactly solvable model results. The label "momentum" is that of the incident γ^* or p.

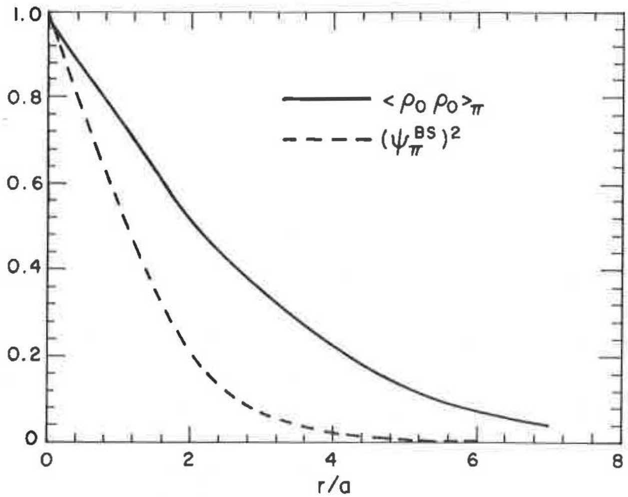

FIGURE 4 Comparison (Ref. 32) of the density–density correlation (solid line) with the square of the Bethe–Salpeter amplitude (dashed line) for the pion. Ψ_π^{BS} measures only the $q\bar{q}$ component of the pion.

That a configuration has a small size might seem strange to quark modelists. The quantum fluctuations of the hadronic size are often large, and one does not characterize the quark wave function by configurations of different sizes. Nevertheless, there is some non-vanishing probability that all of the constituents are close together. Suppose the most important contributions to a matrix element occur from regions in which all of the quarks are close together; then that matrix element measures the properties of a PLC.

The important question concerns high momentum transfer semi-exclusive processes. Do such processes necessarily involve a PLC? We examine simple models to see.

4.1. Nonrelativistic Constituent Models of Hadrons

At high momentum transfer, the nucleon form factor is the matrix element of the electromagnetic current. This hard scattering operator T_H acts on a nucleon to form a wave packet or ejectile. Our interest is in the transverse size distribution of this ejectile because the ejectile–nucleon interaction depends on this size (cf. Eq. (1)).

Thus it seems reasonable to consider nucleonic matrix elements of the operator b^2T_H. Here $\mathbf{b}$ represents the transverse separation between any of the two quarks in the ejectile. The importance of final state interactions is measured by the ratio of the scattering term to the single scattering term so that one wants to consider a ratio defined as $b^2(Q^2)$:

$$b^2(Q^2) \equiv \frac{\langle N(\mathbf{q})|b^2T_H|N(\mathbf{0})\rangle}{\langle N(\mathbf{q})|T_H|N(\mathbf{0})\rangle}\,. \tag{20}$$

The above notation specifies that a nucleon at rest absorbs a momentum $\mathbf{q}$, and the denominator is $F(Q^2)$. Spin effects are again ignored. If the ejectile is a PLC, then the operator b^2 takes on small values and $b^2(Q^2)$ is small. The vanishing of b^2 is a necessary condition for the occurrence of color transparency, but it is not sufficient. One also needs high enough energy such that the escape time is small compared to the expansion time (cf. Section 3). The numerator of Eq. (20) occurs in Eq. (11) if $p_m \approx p$ and Eq. (1) is used.

Consider first a non-relativistic treatment. To gain intuition, start with a two quark object. Thus the elastic form factor is $F(Q^2)$ with

$$F(Q^2) = \int d^3r\psi^*(r)e^{i\mathbf{q}\cdot\mathbf{r}/2}\psi(r) \tag{21}$$

where $Q^2 = \mathbf{q}\cdot\mathbf{q}$, $\mathbf{q} = q\hat{z}$, $Q = q$, $\mathbf{r} = z\hat{z} + \mathbf{b}$ and $\psi(r)$ is the coordinate space representation of the wave function $|N\rangle$. Here $T_H = e^{iqz/2}$, so $b^2(Q^2)$ is given by

$$b^2(Q^2)F(Q^2) \equiv \int d^3r\psi^*(r)b^2e^{iqz/2}\psi(r). \tag{22}$$

One may use $b^2 = r^2 - z^2 \rightarrow -4(\nabla_q^2 - \partial^2/\partial q^2)$ acting on $F(Q^2)$ of Eq. (21) to show that

$$b^2(Q^2)F(Q^2) = -\frac{8}{Q}\frac{dF(Q^2)}{dQ}. \tag{23}$$

So we are now ready to compute $b^2(Q^2)$ for simple models. Start with the oscillator model wave function $\psi^2(r) = ce^{-\alpha r^2}$. Then $F(Q^2) = e^{-Q^2/16\alpha}$ and $b^2(Q^2) = 1/\alpha$. There is no color transparency in this model, because the wave function is a product of functions of z^2 and b^2. This feature is peculiar to the ground state of the harmonic oscillator.

The Coulomb potential provides a different example. Then $\psi^2(r) = ce^{-\alpha r}$, and Eq. (23) yields $b^2(Q^2) = 8/[(Q/2)^2 + \alpha^2]$. Thus color transparency occurs if $Q^2/4\alpha^2 >> 1$.

Both a confining force and a one gluon exchange Coulomb force are important features of quark models. Thus we consider the potential:

$$V = \frac{\beta_1}{2}r^2 + \frac{\beta_2}{r}.$$

It is realistic to assume values $\beta_1/2$ and β_2 such that the term $\beta_1 r^2$ dominates in the calculation of energies of eigenstates, $\beta_2 < 0.1(\beta_1/m^3)^{1/4}$. A calculation of the form factor shows that b^2 is independent of Q^2 at low Q^2, but that at large Q^2 the second term in V dominates. Indeed, for $QR \gtrsim 1$,

$$F(Q) \approx \frac{\beta_2}{\sqrt{\pi}} (2/QR)^4$$

according to numerical calculations. Here $R = (m\beta_1)^{-1/4}$ is the harmonic oscillator length parameter and m is the constituent quark mass. This form factor also gives $b(Q^2) \sim 1/Q^2$.

The above examples indicate that the presence of a $1/r$ potential leads to a small value of $b^2(Q^2)$, for not too high values of Q^2. Thus small values of $b^2(Q^2)$ are natural in constituent quark models.

The general condition needed to obtain a Fourier transform with a power law falloff can be obtained from asymptotic expansions. One finds that $F(Q^2) \sim 1/Q^4$, if $d\psi^2/dr$ does not vanish at the origin. This derivative vanishes only if the potential is a continuous analytic function of r^2.

Next we consider systems of three quarks. The relevant version of the operator b^2T_H is

$$b^2T_H(Q^2) = \sum_{i<j} (b_i - b_j)^2 \sum_k e^{iqz_k} e_k. \tag{24}$$

Here e_k is the electric charge of the k'th quark. Knowledge of the form factor is not generally sufficient to determine $b^2(Q^2)$ for baryons.

One can relate $b^2(Q^2)$ to $F(Q^2)$ in mean field (MF) models with, e.g., $\langle r_1, r_2, r_3|\Psi\rangle = f(r_1, r_2)g(r_3)$. Then Eq. (24) leads to

$$b^2_{\mathrm{MF}}(Q^2)F(Q^2) = \frac{2}{3}\left(\frac{-2}{Q}\right)\frac{\partial F(Q^2)}{\partial Q} + \frac{1}{3}F(Q^2)\langle\Psi|(b_1 - b_2)^2|\Psi\rangle. \tag{25}$$

The term proportional to the form factor arises from the three quark terms of b^2T_H.

As an example of a MF model, suppose each quark has a constant density within a radius R and the nucleon space wave function is a product of the three quark wave functions. (This is the quark part of the SUNY chiral bag model,[33] for chiral angle $\theta = \pi/2$.) Then $F(Q^2) = 3j_1(QR)/QR$, and

$$b^2(Q^2)/b^2(0) = \frac{2}{3} - \frac{1}{3F(Q^2)}\frac{5}{Q}\frac{dF(Q^2)}{dQ} . \tag{26}$$

Note the presence of a constant term, and that the zeros of $F(Q)$ cause rapid increases in $b^2(Q^2)$.

In modern constituent quark models (e.g., Refs. 34–36), quark–quark interactions are very important in determining the energy and wave function. Thus, we do not expect these models to behave as mean field approximations. To verify this, we compute $b^2(Q^2)$ using the wave function of Capstick and Isgur.[34] The results are shown in Fig. 5. As expected, full color transparency is obtained.

4.2. Models with Pion Clouds

So far the nucleon has been considered as a three-quark system. This is, at best, an oversimplification. In QCD the nucleon is expected to have many Fock space components. One way to partially account for these effects is to include a pion cloud or to treat

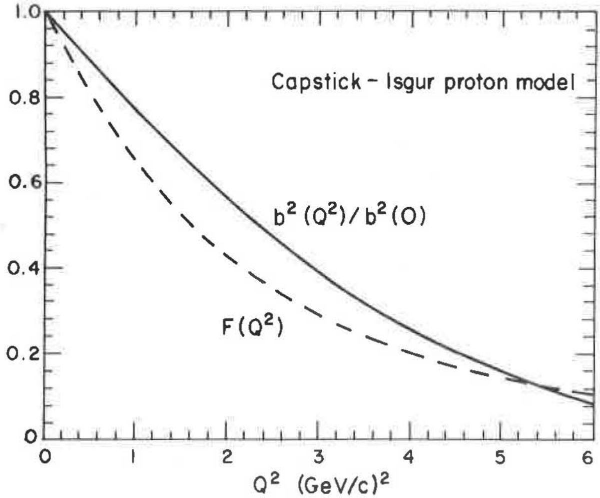

FIGURE 5 Capstick–Isgur proton wave function.

the entire nucleon as a soliton of pion fields. In this Skyrmion case, the baryonic matter density $\rho(r)$ is expressed as

$$\rho(r) = \frac{\sin^2\theta(r)}{r^2}\frac{d\theta}{dr} \tag{27}$$

where **r** is the distance from the bag center, and $\theta(r)$ is a chiral angle representing the radial dependence of the pion field. The Skyrmion model has many desirable features,[37] so it is worthwhile to see whether it possesses a PLC.

To study the role of the PLC in the Skyrme model we define a form factor $F(Q^2)$ and $b^2(Q^2)$ in analogy with Eqs. (21)–(23):

$$F(Q^2) = \int d^3r e^{iqz}\rho(r), \tag{28a}$$

$$b^2(Q^2)F(Q^2) = \int d^3r e^{iqz}b^2\rho(r). \tag{28b}$$

One may again do the angular integral to obtain an expression similar to Eq. (23).

The motivation for the b^2 factor appearing in the integral of Eq. (28b) is the wish to probe the interactions which stabilize the Skyrmion. This interaction may also be influenced by color screening effects. This is possible since at high density the Skyrme model dynamics is expected to be very similar to the behavior of a gas of free quarks.[38] The baryon density is used in Eq. (28) because strong final state interactions are studied. The present study is a simple first investigation. Improvements by readers are encouraged!

We proceed with some numerical evaluations of Eq. (28). The function $\theta(r)$ is taken from the work of Adkins and Nappi[39] which includes the effects of a non-vanishing pion mass. Figure 6 shows that b^2 decreases noticeably with Q^2.

The Stony Brook group[33,40] modified the Skyrmion model to include quarks confined inside an MIT bag of radius R, and a pion soliton outside. The Skyrmion model has $\theta(r = R) \equiv \theta = \pi$ with $B_q = 0$. Conversely if $\theta = 0$, $B_q = 1$ and the MIT bag model limit is obtained. The computed matrix elements are independent

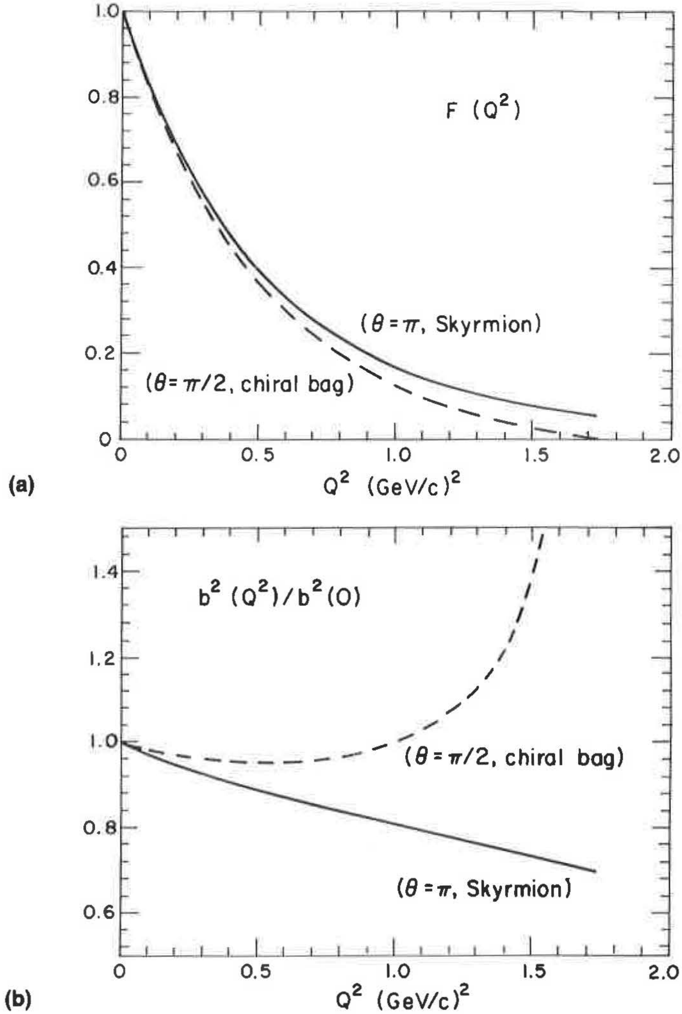

FIGURE 6 Skyrme model (solid) vs. chiral bag (B_q = 1/2). (a) Form factor. (b) $b^2(Q^2)$.

of R to a good approximation. This is the cheshire cat principle.[41,42] However, the $b^2(Q^2)$ are very different even though the form factors are very similar. Results for $\theta = \pi$ and $\theta = \pi/2$ ($B_q = B_\pi = 1/2$) are shown in Fig. 6. This is because $b^2(Q^2)$ is sensitive to higher moments of r than $F(Q^2)$. The mean field and sharp surface nature of the quark density cause the rise shown in Fig. 6b.

Another popular model is the cloudy bag model.[43,44] Here the pion is a quantum fluctuation treated in perturbation theory. The results using the Regensburg version of the cloudy bag[45] are represented in Fig. 7. The Gaussian falloff of the pionic contribution is due to the combination of the Yukawa behavior with the surface peaked nature of the π-quark coupling.

A comparison of the pion and quark contributions to the form factor shows that the pion tail which is relevant for a considerable part of $b^2(Q = 0)$ becomes inessential at $Q \gtrsim 1$ GeV^2. This interesting effect can be studied in the domain of traditional nuclear physics. The final state interaction of a nucleon should be suppressed at $Q^2 \sim 1$ GeV^2 since the effective size is reduced by the loss of the pion cloud. This effect is reduced to some extent by

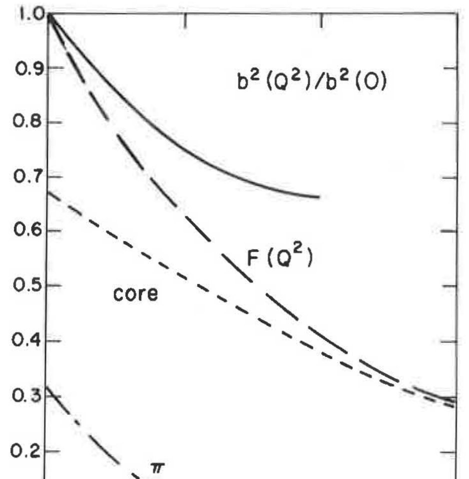

FIGURE 7 Cloudy bag model calculation of Ref. 45. The quark core (dotted line) and pion cloud (dashed line) contributions to the form factor (long dashed line) are shown. $b^2(Q^2)/b^2(0)$ is the solid line.

the quantum diffusion of the PLC (recall Section 3). We estimate the time for the low momentum pion cloud to reappear to be of the order of $1/E_\pi \approx 1/2m_\pi$. The vanishing of the pion cloud can be essential at distances less than $(E_N/m_N)\ 1/2m_\pi \simeq (1 + Q^2/2m_N^2)/2m_\pi \approx 2$ fm if quasielastic kinematics are used and $Q^2 \sim 2$ GeV^2.

The suppression of the pion cloud around the interacting nucleon can be investigated in the reactions $(e, e'N)$, $(e, e'\Delta)$ of the lightest nuclei. For example the quasielastic reactions ($e + A \rightarrow e + \Delta^{++}$, $e + \Delta^-$) dominated by the final state interaction ($\Delta^+ p \rightarrow \Delta^{++} n$) should be strongly suppressed as compared to expectations of the Glauber approach. Indeed, charge exchange two body processes $p + p \rightarrow \Delta^{++} + n$, $n + n \rightarrow \Delta^- + p$ are well described in terms of the one pion exchange diagram in a wide kinematical region.[46] Thus, suppressing the pion leads to suppressing the Δ^{++} or Δ^-. Light targets are necessary for such experiments. This reduces the chance that the Δ decays in the nucleus as well as the probability for the reappearance of the pion cloud. The Q^2 needed is about 2 GeV^2, so the experiment could be done at CEBAF energies.

4.3. Relativistic Effects

The non-relativistic models of the previous sub-section are somewhat suspect, since high momentum transfer is important. A relativistic treatment is therefore mandatory, even though existing methods are not completely satisfactory.

First, we consider the dependence of $b^2(Q^2)$ in terms of the light-cone quantum mechanics of a two-quark system. The form factor is given by

$$F(Q^2) = \int \frac{d\alpha d^2k_t}{\alpha(1 - \alpha)} \psi^*(\alpha, \mathbf{k}_t)\psi(\alpha, \mathbf{k}_t - \alpha\mathbf{q}_t). \qquad (29)$$

Here $q = (q_+, q_-, \mathbf{q}_t)$. We choose $q_- = q^0 - q^3 = 0$, $q^2 = -Q^2 = -q_t^2$ and the direction of $\mathbf{q}_t$ as the x axis. This light-cone physics formula[47,48] describes the absorption of a photon by one of the constituents. The quantity α is the fraction of hadron plus momentum carried by the spectator.

The expression (29) can be used to gain qualitative insight about the transverse sizes involved at high momentum transfer process. Suppose the integral is dominated by contributions from regions where α is small. Then k_t takes on values controlled by the size of the wave function and small transverse sizes are not important. In this case, the struck quark has the high momentum fraction $1 - \alpha$. So one sees the Feynman mechanism (recall Section 1.2). On the other hand, if α is not a very small fraction, the relevant values of k_t must be of the order of q_t. Then small transverse sizes must be involved for sufficiently large values of q_t.

A relativistic version of $b^2(Q^2)$ is needed. To obtain such, we express Eq. (29) in terms of the two-dimensional Fourier transform $\Phi(\alpha, b)$ of $\psi(\alpha, k_t)$. Then a relativistic version of $b^2(Q^2)$ is defined in analogy with Eq. (8):

$$b^2(Q^2)F(Q^2) \equiv \frac{1}{2}\int \frac{d\alpha d^2b}{\alpha(1-\alpha)}\,\Phi^*(b,\alpha)b_y^2 e^{iq_t b_x \alpha}\Phi(b,\alpha). \quad (30a)$$

The matrix element of b_y^2 is used because we are concerned with the distribution of constituents in the direction transverse to $\mathbf{q}$. The normalization is chosen so that in the non-relativistic limit (here $\alpha \approx 1/2$) Eq. (30a) leads to the usual result that $b^2(Q^2 = 0)$ is 2/3 of the mean square radius.

Equation (30a) can be rewritten in the momentum space representation. The result is

$$b^2(Q^2)F(Q^2) =$$

$$-\frac{1}{2q_t}\frac{\partial}{\partial q_t}\int \frac{d\alpha}{\alpha(1-\alpha)}\frac{d^2k_t}{\alpha^2}\,\psi(\alpha,\mathbf{k}_t)\psi(\alpha,\mathbf{k}_t - \alpha\mathbf{q}_t). \quad (30b)$$

Equation (30) and its non-relativistic version Eq. (24) are the principle new results of this paper.

Our first relativistic example is "the harmonic oscillator,"

$$\psi^2(\alpha, k_t) = Ae^{-B[(m^2+k_T^2)/\alpha(1-\alpha)-4m^2]}, \quad (31)$$

where m is the quark mass. For large values of q_t the integrals (29) and (30) can be performed analytically for k_t and with the method of steepest descent for α. The saddle point occurs for small values of α: $\alpha \approx 2m/q_t$. Thus, the harmonic oscillator provides an example of the Feynman mechanism. At high $q_t = Q$, $F \sim (1/Q^{3/2})e^{-BmQ}$, see Ref. 49. The use of Eq. (30) gives $b^2(Q^2)F(Q^2) \sim -(\partial/Q\partial Q)(Q/2m)^2F(Q^2)$ so that $b^2(Q^2)/b^2(Q^2 = 0) \sim Q/m$. The relativistic harmonic oscillator wave function with its Feynman mechanism (end point singularity) leads to color opacity. Numerical results are shown in Fig. 8. The quark mass is 330 MeV and the rms radius is 0.8 fm.

Note that using a wave function to reproduce an observed form factor is not sufficient to prove or disprove the existence of a PLC. Consider a "hydrogen" wave function:

$$\psi_h(\alpha, k_t) = \frac{A}{\left[\dfrac{m_q^2 + k_t^2}{\alpha(1 - \alpha)} + \mu_h^2\right]^2}. \quad (32)$$

The evaluation of $b^2(Q^2)$ now leads to color transparency at $Q^2 \geq 1$ GeV^2 (Fig. 8). We take the same quark mass and rms radius as for the harmonic oscillator.

The "harmonic oscillator" and "hydrogen atom" wave functions and the Capstick–Isgur model are three examples of "soft wave functions" in the sense of Isgur and Llewellyn-Smith[50]: no explicit gluons are present. The first wave function provides color opacity while the other two yield transparency. Thus a soft wave function can yield color transparency. The validity of pQCD is not a requirement for the existence of color transparency.

The analyses of Sections 4.1–4.3 show that existing hadronic models provide very diverse predictions for the emergence of color transparency. It occurs naturally in realistic quark models where correlations caused by gluons are important, in the Skyrmion model and in lattice calculations. At the same time, the QCD sum rule theory of hadronic form factors[8] predicts no color transparency at all accessible momentum transfers $Q^2 < 10$ GeV^2; while the relativistic harmonic oscillator model leads to color opacity. There-

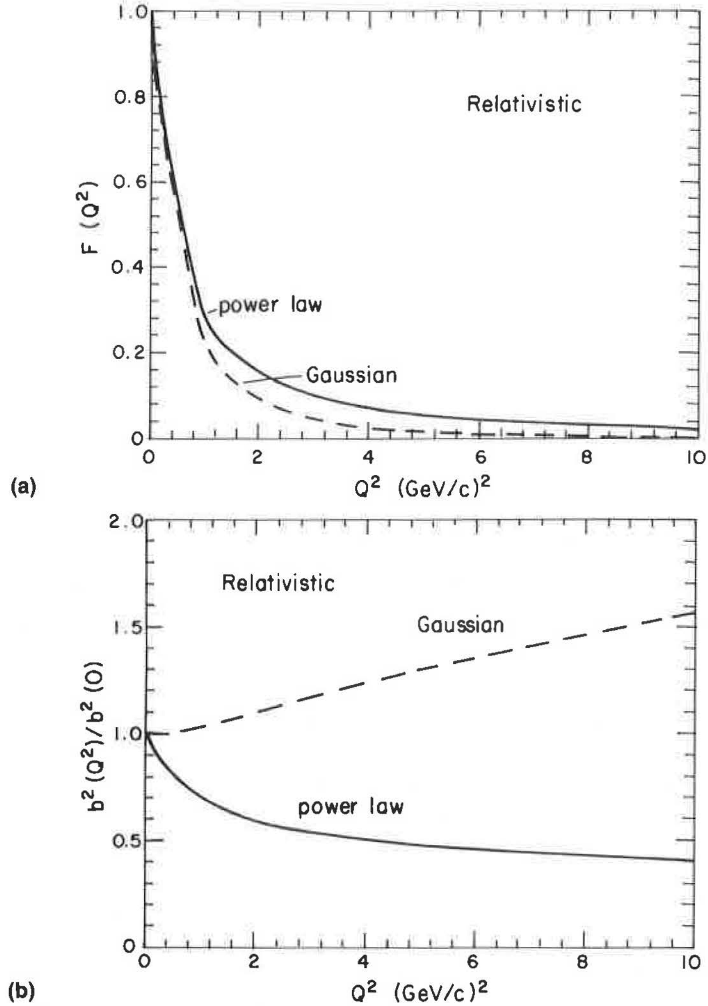

FIGURE 8 Relativistic harmonic oscillator (dashed) and hydrogen atom (solid). (a) $F(Q^2)$. (b) $b^2(Q^2)/b^2(0)$.

fore consituent quark models and QCD sum rules can be distinguished by whether or not color transparency is observed at these Q^2. We also find that the observation of color transparency would not necessarily imply the validity of pQCD calculations of hadron form factors.

4.4. PLC and the Nuclear Ground State

So far, we have discussed whether a PLC can be formed in a high momentum transfer process. However, the presence of such a component in a bound nucleon could have important implications for ground state nuclear properties.[1,21] The form factor of a bound nucleon should be different from a free one. This could occur via the suppression of the small momentum Fourier component of the pion field of a bound nucleon due to the Pauli principle[51] along with the suppression of interaction with other nucleons of bound nucleon in the PLC configuration. As a result the probability of PLC is suppressed in a bound nucleon by the factor $\delta(\kappa)$.[21,23] Here κ is the momentum of the bound nucleon in the nucleus and

$$\frac{F^{\text{bound}}(q^2)}{F^{\text{free}}(q^2)} = \delta^{1/2}(\kappa) = \left\{1 + \frac{2(\kappa^2/2m) + \epsilon_A}{\Delta E}\right\}^{-1/2} . \quad (33)$$

The suppression of Eq. (33) takes place at small κ where the $N\,N$ interaction is soft. At large κ, PLC are likely to dominate in the $N\,N$ interaction and $\delta(\kappa) > 1$.

This physics may reveal itself in the difference between form factors of free and bound nucleons at large nucleon momenta in $(e, e'p)$ and $(h, h'p)$ reactions. It may be also responsible for the $x > 0.3$ EMC effect.[21] If so, $\Delta E \sim 0.6$ GeV.

4.5. Physics of Heavy Ion Collisions

The fluctuations of color in a nucleon lead to a number of interesting consequences for the physics of central heavy ion collisions at CERN nuclear beam energies (200 GeV/nucleon).

1. The production of the leading (highest energy) nucleons and nucleon isobars should rapidly increase with initial beam energy becoming noticeable at CERN nuclear beam energies (200 GeV/nucleon). This is because more PLC can be considered as frozen during the scattering process.[52]

2. Large fluctuations of projectile size lead to large variations in multiplicity and transverse energy from event to event.[53] These fluctuations are the consequences of diffractive processes that are

important. Thus a dedicated investigation of diffractive processes at RHIC energies is essential.

The investigation of color transparency found a rough equivalence between the description of the process in terms of hadronic intermediate states and in terms of fluctuations of the space size of a configuration (and cross section) depending on the distance from the point where particle (jet) is produced. It is interesting that even at not large energies Skyrmion model calculations found a similar equivalence. The radius of a Skyrmion vibrates as a function of distance from the point of collision.[54] Thus it is possible to improve calculations of cascades by including quantum mechanical effects due to inelastic intermediate states. This physics can be modelled by using the $\sigma(Z)$ of color transparency phenomena.

5. DISCUSSION OF AVAILABLE DATA

The single published experiment aimed at observing the effects of color transparency is that of Carroll *et al.*[55] This was the BNL (p, pp) work at beam momenta p_L ranging from 6 to 12 GeV/c. In (p, pp) experiments the projectile and struck proton must form a two baryon PLC if color transparency is to be obtained. The kinematics of the BNL experiment were set so that the basic pp elastic scattering event occurs at a center of mass angle of 90° if the target proton is at rest. The data for the transparency ratio are shown in Fig. 9. The data presented include many theoretical corrections. One example is the effect of Fermi motion which is used to obtain the points at momenta different from p_L. Obviously this involves a model for the nuclear wave function. Another example is lack of a direct measurement of the excitation energy of the recoil system—for this one needs to measure the momenta of the two outgoing nucleons (in Ref. 55 only the emission angle ($\hat{p}$) was measured for one of the emitted nucleons, and information from veto counters was used to suppress inelastic events). There is also a problem of calculating off-shell corrections of the order m^2/s which would be negligible at large energies but not in most of kinematics studies in Ref. 55. These arise because the value of s for the projectile–target proton system need not be the same as

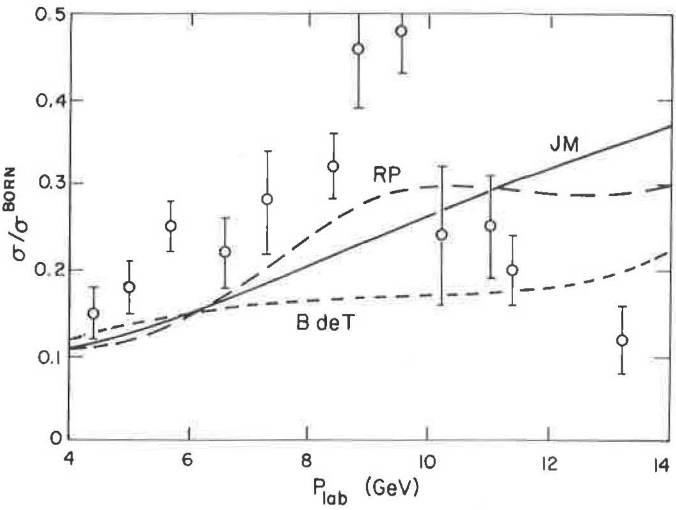

FIGURE 9 $^{27}Al(p, 2p)$. The data are from Ref. 55. The curves are JM (Ref. 10), RP (Ref. 47) and B deT (Ref. 62).

for the two outgoing protons. In any case, we discuss the data as they exist.

Figure 9 shows a transmission T (ratio of nuclear to hydrogen cross section per nucleon after removing the effects of nucleon motion) with an oscillatory pattern. T increases as the beam momentum increases from 5 GeV/c but then decreases. The decrease was unexpected and there is still no unanimously accepted explanation for this. Also shown is the expectation[5,10,55] based on standard Glauber theory that does not reproduce the data. The standard survived a rigorous examination in Ref. 56. The Glauber curve is always independent of energy, but the magnitude depends on whether the proton total or reaction cross section is used.

One possibility, suggested by Ralston and Pire,[57] is that the transparency oscillates as a function of energy due to the interference between a hard amplitude which produces a small object and a soft one (the Landshoff process[58]) which does not. Quark interchange diagrams must also be included to avoid a contradiction with the experimental observation that the large angle pp and

$p\bar{p}$ cross sections differ by a factor of 100 ($pp > p\bar{p}$). At high energies this mechanism leads to approximately the same Q^2 dependence as hard scattering.[59] At present energies, the energy dependence and phase of the Landshoff term are not known. Moreover, it is not even clear if the Landshoff process produces a large object,[60] and the leading log asymptotic perturbative QCD calculations find a negligible phase difference between the two amplitudes.[60]

Zakharov and Kopeliovich[20] and Jennings and Miller[61] have pursued the Ralston–Pire idea with the aim of including the effects of PLC expansion. More careful treatments of the nuclear interactions of ordinary components of the proton were made and the effects of quantum diffusion were included. Some results are shown in Fig. 9, and the data are not well reproduced.

Another idea[62] is that the two-baryon system couples to charmed quarks and the transparency decreases as a result of opening up $\bar{c}c$ channels at energies near the threshold. This is motivated by the fact that the mass scale of the rapid energy variation[63] in A_{NN} and in the measured transparency matches that of the charm threshold. Thus, in this threshold model the drop in color transparency is tied to the increase in A_{NN}. The color transparency predicted by the Brodsky de Teramond idea had never been evaluated until the work of Ref. 60. Their results are also shown in Fig. 9. Once again the data are not well reproduced.

Our investigation of models finds a variety of effects. There is transparency at moderate Q^2 due to the disappearance of the pion field, and oscillations between N and N^* as a function of atomic number and energy. But opacity occurs in the mean field approximation, and transparency occurs for still larger Q^2 as a result of short range (gluon exchange) interactions between quarks (Section 4). So there is much more to be learned. The new EVA experiment[64] designed for higher energies and greater accuracy will certainly help.

We explained above that expansion of the wave packet tends to mask effects of CT at intermediate energies in large angle reactions. To suppress this effect it is necessary to consider reactions where the rescattering cross section in the hard process is measured directly and the nucleons of the projectile involved in the process are sufficiently close. The simplest examples are reactions[1]

$$e + D(^4\text{He}) \rightarrow e + p + n(p + p + (nn)), \qquad (34a)$$

$$h + D(^4\text{He}) \rightarrow h' + p + n(p + p + (nn)) \qquad (34b)$$

in the kinematics where Q^2, $t = -(E_h - E_{h'})^2$, $t_i = -(p_D - p_{N_i})^2$ are large, and $(p_{N_1} + p_{N_2})^2 >> 4m_N^2$.

In these kinematics (Fig. 10) the struck nucleon in reaction (34b) has to pass only 1–1.5 fm before the second rescattering—therefore the expansion effects are of much less importance. If we fix t corresponding to $\theta_{c.m.} \sim 90°$ and select t_1 large enough, we can study the rescattering probability $P(t_2)$ as a function of t_2. For small $t_2 P(t_2) < P_{\text{Glauber}}$ due to CT. But at large t_2 probability of rescattering is enhanced.

$$P(t_2)/P_{\text{Glauber}} \sim 1/F_N^2(t_2) \sim (t_2)^2$$

since only the PLC contributes in the process. One does not have to pay twice the price of finding the projectile in the PLC.

We also show predictions for the $(e, e'p)$ reactions. The kinematics are of the quasielastic type, with the energy dependence of computed ratios of cross sections of ^{12}C (Fig. 11). The curves are from Farrar *et al.*[5] and Jennings and Miller[10] who both made the perhaps optimistic assumption that a PLC is formed in the elementary high momentum transfer process. This is in line with our analysis of realistic models of a nucleon.[34–36] The quantities σ are $(e, e'p)$ differential cross sections integrated over the scattering

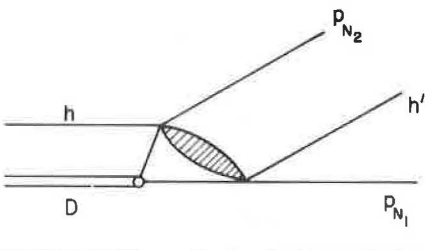

FIGURE 10 Kinematics for double scattering.

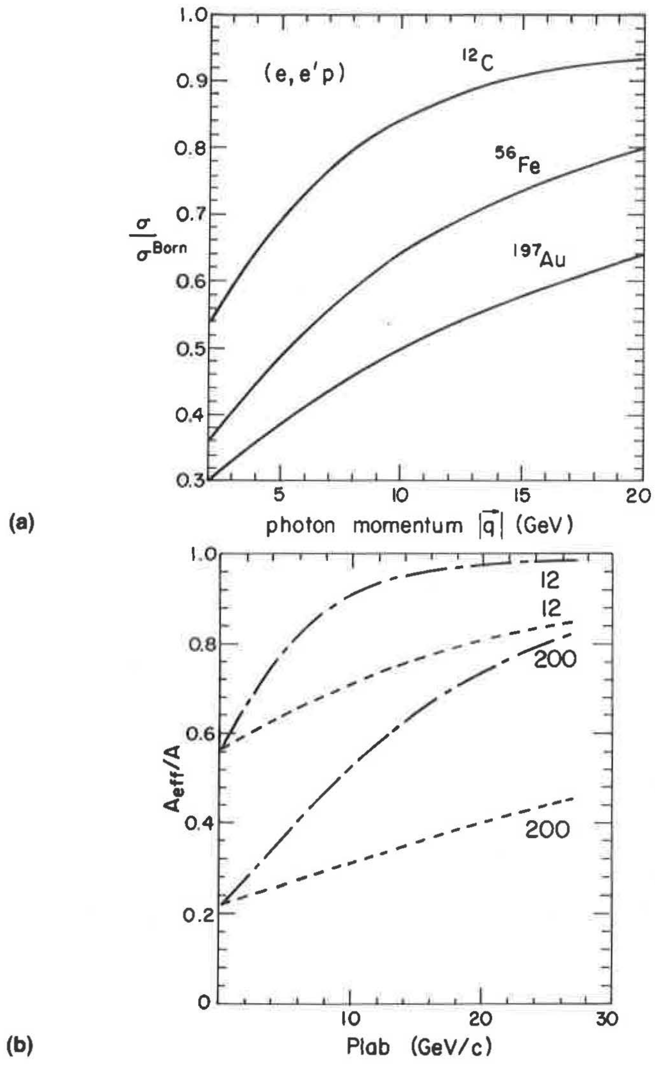

FIGURE 11 (*e*, *e′p*). (a) Predictions of Ref. 10. (b) Predictions of Ref. 5.

angles of the outgoing proton. Full color transparency corresponds to unity. We are concerned with the energies for which the ratio approaches unity and the energies for which the ratio is substantially greater than that obtained with the standard Glauber treatment. We see from Fig. 11 that observable increases are obtained for values of **q** as low as 5 GeV/c, or $Q^2 = 9$ GeV/c^2. Furthermore the strength of such enhancements depends critically on the time scale parameter (recall Section 3).

These predictions help experimentalists to realize that color transparency is not an academic exercise. An experiment at the SLAC NPAS facility[65] has run and is in the process of being analyzed. The present experiment covered the accessible range from $Q^2 = 1$ to $Q^2 = 7$ GeV^2 and has found that the extension to values ≤ 15 GeV^2 is possible despite the low duty factor of the SLAC linac. A noticeable enhancement can be seen. The reader will note, however, that we show a ratio. The numerator is to be measured, but the denominator can't be measured. One must rely on calculations that depend on nuclear wave functions. Only the well-known momentum transfer aspects of these wave functions are needed here. But high accuracy is important. The best way to reduce these kinds of uncertainties is to study the energy and A dependence of the ratios. It is clear that measurements at higher values of Q^2 are also very important.

Another interesting question concerns the transverse sizes important in large t (e.g., $-t > 2$ GeV^2), small angle scattering at large energies. Different QCD diagrams are likely to be responsible for large and small angle scattering, so an experimental study is of interest. One option is the familiar $(p, 2p)$ reaction. Since the incoming particle is fast, its scattering state can be considered as frozen. Therefore if the transverse size decreases as t increases, one should observe that the cross section depends on A as $A^{2/3}$ instead of $A^{1/3}$. This is due to the contribution of scattering from the back surface of the nucleus. In other elastic scatterings, the recoiling proton is likely to be absorbed, since it expands rapidly after the collision. At $-t > 10$ GeV^2, the cross section may start to approach A since the recoiling nucleon can propagate distances > 2 fm in the PLC.

6. SUMMARY AND PERSPECTIVES FOR THE FUTURE

The physics of color transparency CT involves a full intersection of particle and nuclear physics ideas and techniques. It may become an effective method to investigate the physics of confinement and spontaneously broken chiral symmetry in QCD. Such issues are important aspects of heavy ion collisions, so that CT physics may help there, too.

We find (Section 4) that realistic models of a nucleon do allow the formation of a point-like configuration PLC (one of the requirements for color transparency to occur) in high momentum transfer quasielastic reactions. PLC occur in hadrons as described by realistic quark models, in the Skyrmion model, and in lattice calculations. In models with pion clouds, the physical nucleon can become smaller because of the disappearance of that cloud.

However, the QCD sum rule result of Ref. 8 is that no PLC are made at all accessible momentum transfers $Q^2 \leq 10\ \text{GeV}^2$. Thus, observing, or failing to observe, CT effects can rule out theories of hadronic form factors. In addition, the validity of pQCD for explaining measured hadronic form factors is not a requirement for color transparency (Section 4.3).

We have discussed (Section 3) that projectile-size fluctuations can be described in terms of a wave packet of hadronic intermediate states. An exactly soluble model of these fluctuations was presented. In simplified situations this model yields predictions of oscillations similar to those observed in Ref. 55. (This model may be more suited to the study of the nuclear production of $c\bar{c}$ states.)

Studying color transparency is a theorist's delight. Many novel effects can be predicted. But we stress that computing color transparency effects is not simply a theoretical exercise. Interesting effects may occur in the kinematical region accessible (Q^2 as low as 1 GeV^2) for current and planned accelerators. So dedicated experimental and theoretical investigations in this rapidly developing field are necessary.

We outline some possible future directions and unresolved problems, starting with experimental issues.

1. Electron, proton, pion and kaon beams available or to be available at SLAC, BNL, KAON and at the upgraded FNAL fixed target beam lines can be used to search for color transparency effects. Such measurements would help to resolve the long-standing question of whether wide angle two-body reactions at feasible energies are dominated by small or large interquark

distances. The PLC can be described as a coherent sum of baryon excitations. Thus, dedicated measurements of baryon resonance production are urgently needed.

2. Transparency and the disappearance of the pion field in moderately hard ($Q^2 \sim 2$ GeV^2) collisions are phenomena unexpected in traditional low energy physics. Thus, observing such effects would have a strong impact for constructing bridges between nuclear physics and QCD. An experiment necessary to see the disappearance of the pion field would involve the suppression of the quasielastic formation of resonances which cannot be produced in two-body processes. The production of Δ^{++} and Δ^- by electrons and in $A(p, p\Delta)$ reactions are only two examples. The reader is encouraged to find other cases in which CT is manifest by the suppression of a nuclear reaction. Thus CT predicts the enhancement of some reactions and the suppression of others.

Now we turn to future theoretical investigations.

1. The PLC can be described in terms of a bare quark–gluon Fock space. Those configurations with the minimal number of bare particles in hadronic wave functions are candidates to be a PLC. This Fock space description is simple and applicable to many different processes: elastic form factors at high Q^2, deep inelastic scattering at $x \sim 1$, and large angle Compton scattering. Such configurations may be studied with lattice techniques, so there is a possibility to test QCD in a new way. Thus, the pressing problem is to develop methods to calculate wave functions of the $3q(q\bar{q})$ configuration in a nucleon (meson).
2. Evidence that the nucleon radius vibrates has been obtained in calculations of Skyrmion–Skyrmion scattering. For collision kinetic energies ~ 0.5 GeV, the Skyrmions are small at small separations (see Fig. 7 of Ref. 38). It seems fruitful to investigate transparency by studying a third Skyrmion in the presence of a PLC formed by two others.
3. The hadronic size fluctuations of CT phenomena are also relevant for constructing theories of phase transitions in superdense nuclear matter and heavy ion collisions. For example, cascade codes can be modified to include the dependence of

the cross section on the hadron size and on the distance from the production point. This would allow quantum mechanical effects to be included in classical mechanics calculations. Another example is to account for non-nucleonic degrees of freedom in constructing theories of superdense nuclear matter (e.g., the neutron star core). One can use equations like (14) to calculate the admixture of baryon resonances.

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