Fluctuating Proton Size and Oscillating Color Transparency

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The unexpected energy dependence of high-energy fixed-angle *pp* elastic scattering in a nuclear target can be interpreted in terms of interference between two perturbative QCD subprocesses. By proposing the attenuation of Landshoff-type contributions in nuclei, we obtain a parameter-free relation that matches the energy dependence of the data of Carroll *et al.* The approximations improve with larger A and higher energy, leading to a prediction of oscillatory energy dependence for the transparency ratio at higher energy.

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In QCD the proton is viewed as having many components in its Fock-space wave function, from three quarks to an almost unlimited number of soft charged partons and gluons. If many modes are superposed to make a physical proton, the amplitude in any particular mode, e.g., three quarks, must show time-dependent fluctuations. The idea that the proton contains a significant amplitude to fluctuate to a spatially small, color-neutralized configuration has led to the prediction[1] that nuclear matter might be anomalously transparent to protons undergoing elastic scattering through large momentum transfer. To test this, the cross section $d\sigma/dt$ for fixed-angle elastic scattering at 90° c.m. has recently been measured[2] in different nuclei by Carroll *et al.*

The key ingredient in the transparency arguments is that protons that undergo wide-angle, high-energy scattering should be small. In perturbation theory, quark-counting diagrams[3] confirm this, as each scattered quark propagator is someplace off shell by order $1/s$. However, there is a class of graphs, pointed out by Landshoff,[4] in which independent hard scatterings occur with no far-off-shell quarks. For these "disconnected" graphs, one cannot show that the hard-scattered proton is small. Indeed, the amplitude for such a process *increases* if more impact parameter space is available, i.e., if the protons are "large."

We can see this generally from the kinematics. A typical $3+3\rightarrow 3+3$ quark-counting amplitudes scales like $1/s^4$, giving the fixed-angle $d\sigma/dt \propto 1/s^{10}$. In comparison, if $\mu^2 \lesssim p_T^2$ is an internal transverse-momentum scale in the proton wave function, $\mu^2 \ll s$, the disconnected amplitudes scale like $1/(\mu^2 s^3)$, and $d\sigma/dt \propto 1/s^8$. The sensitivity to the p_T scale μ is a relic of the spatial separation of quark collisions in the disconnected scatterings. A fluctuation in the proton size, in an elastic collision, can be thought of as a fluctuation in the effective value of μ. As $\mu \rightarrow 0$, the transverse coordinate space available for collisions of quarks increases and the amplitude is larger. In the extreme limit $\mu \rightarrow 0$, one can visualize independent scatterings of plane waves occurring on opposite sides of the "lab," in fact. In practice, the spatial separation is certainly limited by the size of the proton, of order 1 fm. For any fixed μ, the s^{-8} energy dependence favors the independent scattering processes at high energy over the s^{-10} quark-counting process.

In QCD it is unlikely to find either a "small" proton or to find perfect elastic collisions of the independent hard scatterings. The Landshoff terms are believed to be partially suppressed[5] by soft-gluon contributions related to the QCD Sudakov form factor, which express the (small) amplitude to find a free charge in a gauge theory. Some analysis has been done[6] to show that the suppression due to soft gluons creates a power intermediate between the quark-counting and Landshoff values. Thus the independent diagrams of Landshoff actually dominate at high energy. The intermediate value of the power indicates that the dominant scale for fluctuating μ values does not grow as fast as s.

Further evidence for this picture comes from our considering the interference between the two processes.[7] In the interference, the Sudakov effects become visible through an energy-dependent phase[7,8] which can be calculated in QCD. The phase is due to initial- and final-state interactions as color is separated, creating the QCD version of the Coulomb phase shift. The data for *pp* elastic scattering show a striking oscillatory energy dependence with a logarithmic period that beautifully confirms the picture. In addition, polarization data indicate that large, energy-independent phase differences may have been observed,[9,10] as predicted by the model.[10]

Let us apply the same ideas to the nuclear-transparency experiment. It is clear that large protons will quickly dissipate in a nuclear medium. For an estimate, we use $\sigma_{pp}=40$ mb, $n\sim\frac{1}{6}$ fm^{-3} as the nuclear density, and $L=1/n\sigma\sim\frac{3}{2}$ fm as the mean free path. We compare this to the distance a proton would have to travel, $X(A)$. For a uniform density spherical drop of A nucleons, averaging over impact parameters, then $X(A)$

=radius of drop $=1.2A^{1/3}$ fm. Thus $L \lesssim X(A)$ for practically all nuclei, and large protons are almost certain to collide by conventional soft interactions.

This effect will deplete the Landshoff independent-hard-scattering process. To estimate this, an exponential attenuation is reasonable. Thus if M_L were measured for a proton, its magnitude should change in nuclear targets by about

$$|M_L| \to |M_L| \exp\{-[X(A)-X(1)]/2L\}. \qquad (1)$$

There is a factor of $\frac{1}{2}$ in the exponents in this primitive estimate of the amplitude, not the number flux. Note we are assuming that typical Landshoff-process protons are full sized, i.e., of size $\mu \cong 1/(40 \text{ mb})^{1/2} \cong \frac{1}{2}$ fm^{-1}, well into the "soft" region of QCD. For $A=12$, say, the Landshoff amplitude is depleted by $\exp(-0.5) \sim 0.6$, while for $A=64$ the same amplitude is smaller by about 0.3. We will use this estimate only for an order of magnitude calculation. We note in passing that the estimate is equivalent to the "$\tau=0$" case considered by Farrar, Liu, Frankfurt, and Strikman[11] for the component of scattering with no reduction in size. For the case of small $A \lesssim 10$, the details of attenuation are important and the various models proposed in Ref. 11 can be explored. However, it is important that our predictions do not depend strongly on the details of attenuation, so long as attenuation is significant. We concentrate on the limit $A \gg 1$, where the calculations are more reliable.

In QCD the effects of the energy-dependent "chromo-Coulomb" phase shift for $pp \to pp$ scattering can be represented by the formula[7]

$$M = M_{\text{QC}} + e^{i\phi(s)+i\delta_1}|M_L|; \qquad (2)$$

$$d\sigma/dt = |M|^2/64\pi s p_{1,\text{c.m.}}^2. \qquad (3)$$

Here δ_1 is an uncalculable, energy-independent phase. The calculable phase $\phi(s)$ has a known energy dependence[7,8] analogous to renormalization-group evolution:

$$\phi(s) = \frac{\pi}{0.06} \ln\ln(s/\Lambda_{\text{QCD}}^2). \qquad (4)$$

The quark-counting amplitude is not infrared sensitive, resulting in a phase that is much more slowly varying. Thus (4) is the most important contribution to the energy dependence of the phase difference. The superposition of M_L and M_{QC}, including $\phi(s)$, will result in a term going like $\cos[(\pi/0.06)\ln\ln(s/\Lambda_{\text{QCD}}^2)]$ in the cross section. Thus from Eqs. (2) and (3) we derive $R_1(s)$, the ratio of the cross section to the quark-counting prediction,[7]

$$R_1(s) = s^{10}\frac{d\sigma}{dt}(pp)|_{90°} \propto 1 + \rho_1(s/1\text{ GeV}^2)^{1-K}\cos[\phi(s)+\delta_1] + \rho_1^2(s/1\text{ GeV}^2)^{2-2K}/4. \qquad (5)$$

The constant ρ_1, measuring the relative normalization of M_L to M_{QC}, equals 0.08 from a fit to the region 10 GeV$^2 \lesssim s \lesssim 40$ GeV^2 [Fig. 1(a)]. This is large enough to lead to oscillations of more than a factor of 2 in the data.[7,12] The intermediate power $K=\frac{1}{2}$ and the constant $\pi/0.06$, while consistent with QCD estimates, are chosen by fit with $\Lambda_{\text{QCD}} = 100$ MeV.

For nuclear targets, we propose the relative normalization $\rho_1 \to \rho_A = \rho_1 \exp\{-[X(A)-X(1)]/2L\}$ according to our discussion above. In such a formula we implicitly assume that there is little significant attenuation of the quark-counting amplitudes: they are indeed small. However, only the relative normalization of the two amplitudes is meaningful and it does not seem practical to model in detail the small attenuation of the small protons. In addition, $\delta_1 \to \delta_A$ represents nuclear effects of A dependence on the phase difference. We have no way to estimate these reliably in the procedure we are presenting.

However, if our picture is correct, we should not need the nuclear information for large enough A. For in the case of large A, all large Landshoff protons will be depleted, and their phase shifts $\phi(s)$ and δ_A will become undetectable. For the transparency $T(s)$, defined by

$$T(s) = \frac{1}{A}\frac{d\sigma(pA \to pp(A-1))/dt}{d\sigma(pp)/dt}\bigg|_{90°}, \qquad (6)$$

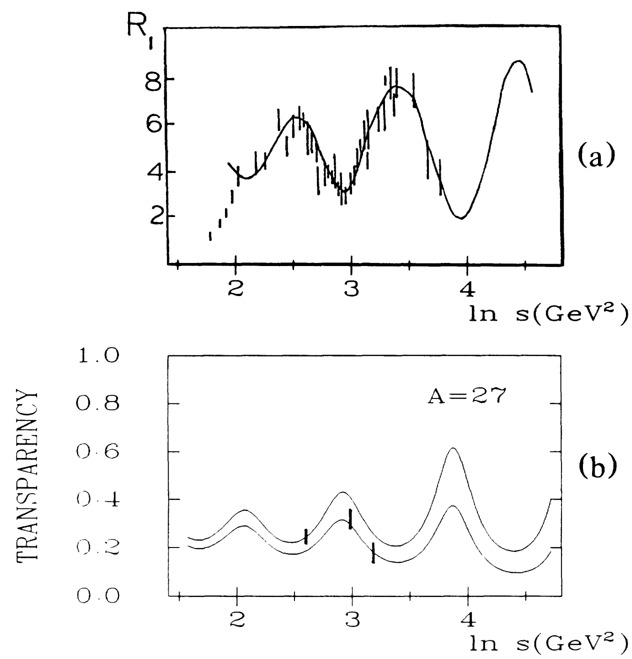

FIG. 1. (a) The energy dependence of $R_1(s) = \text{const}\, s^{10} d\sigma/dt(pp)|_{90°}$ for the high-energy pp elastic scattering at 90° c.m. angle compared to Eq. (5) (solid line), as taken from Ref. 7. (b) Prediction of oscillating transparency $T(s)$ [Eq. (7)] for $A=27$ after varying over all possible nuclear phases δ_A (upper and lower limits are shown); data from Carroll *et al.* (Ref. 2).

we find

$$T(s) = N_A \frac{1+\rho_A s^{1-K}\cos[\phi(s)+\delta_A]+\rho_A^2 s^{2-2K}/4}{1+\rho_1 s^{1-K}\cos[\phi(s)+\delta_1]+\rho_1^2 s^{2-2K}/4}, \quad (7)$$

where N_A is an A-dependent, but energy-independent, normalization. In the limit of large A such that $\rho_A s^{1-K} \ll 1$, the numerator of (7) becomes energy-independent, giving

$$T(s) \to N_A/R_1(s). \quad (8)$$

If these arguments are true, the nuclear medium filters away the Landshoff-type process and destroys the interference phenomena seen in the pp collision. The known energy dependence of $R_1(s)$ of pp scattering can be compared with data for *nuclear* targets.

In Fig. 2 we compare $T(s)$ [Eq. (7)] with data[2,13] for targets with $A \geq 12$. Consistent with our picture that the relative phase of the Landshoff term becomes unimportant, we vary δ_A over all possible values at each energy, resulting in a band with upper and lower limits (dashed lines) on $T(s)$. For large A, e.g., the case of Pb, the calculation is quite insensitive to δ_A, while for small A the bands are comparable to the error bars in size. For reference we also present a solid line calculated with the choice $\delta_A = \delta_1$, to show the effects of $\phi(s)$ in the numerator. It is clear that for all cases the energy dependence is well reproduced with no free parameters. This indicates that transparency in one component of the scattering amplitude has been observed. However, there is no evidence for increasing transparency with energy. Instead, we find strong confirmation for oscillating transparency as the nucleus acts like a filter for large protons. As a prediction, the transparency $T(s)$ should continue to oscillate with energy with geometrically increasing period: this is shown in Fig. 1(b).

The calculation of the A-dependent normalization N_A is a more detailed and model-dependent issue than the energy dependence at fixed A. However, the analysis of Farrar *et al.*[11] shows that N_A should fall with A in several reasonable models. These models and the trend in the data indicate consistency with some important attenuation at current energies, as in our picture.

Brodsky and de Teramond[14] have recently proposed two new $J=S=L=1$ broad dibaryon resonances at $\sqrt{s}=2.55$ and 5.08 GeV to explain the oscillations in the energy dependence of the 90° pp data. By the adjustment of several parameters, the dependence of the double-transverse-spin observable A_{NN} can also be fitted in this way. On the other hand, large values of A_{NN} are also obtained from interference between quark-counting (constituent interchange) and Landshoff triple-gluon-exchange processes such as in our model. Indeed, Brodsky, Carlson, and Lipkin,[15] referring to earlier work of Farrar and Wu,[15] made calculations studying interference between constituent-interchange and Born-term Landshoff models. This showed that A_{NN} could be fitted although Brodsky, Carlson, and Lipkin showed that one could not fit A_{NN} and the $np \to np$ to $pp \to pp$ normalization or the angular distribution simultaneously. But the Brodsky-Carlson-Lipkin and Farrar-Wu analyses were incomplete in lacking the Sudakov corrections which generate the phase $\phi(s)$. Most importantly, the phase $\phi(s)$ is intrinsically a matrix in color space,[7] which rearranges the color weights and quantum number flow in the Farrar-Wu calculation. Incorporating this in the

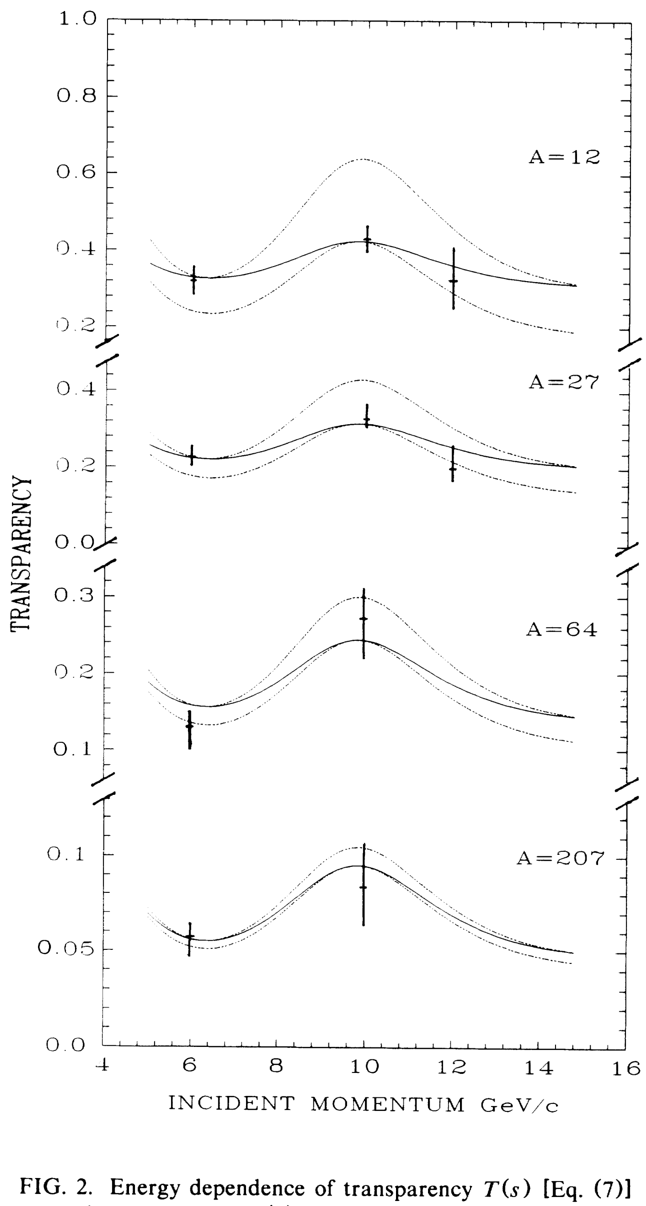

FIG. 2. Energy dependence of transparency $T(s)$ [Eq. (7)] compared to data of Ref. (2) for various A. Upper and lower limits show complete range of sensitivity to nuclear phase δ_A, a range which decreases with increasing A; solid line represents choice $\delta_A = \delta_1$ for comparison.

six-quark-to-six-quark amplitude is a formidable task. Until a complete calculation is done, the question of whether interference between M_L and M_{QC} can explain A_{NN} remains unresolved.[16]

We return to the broader picture of nuclear filtering that we have addressed in this paper. The s dependence of the transparency offers confirmation of interference between two perturbative QCD subprocess we associate with fluctuating proton size. We have not gone so far as to completely specify the nature of the small quark-counting amplitudes in nuclei, but we are able to reproduce the data without need for free parameters by simply attenuating the large amplitudes responsible for energy-dependent phase shifts. Further tests of the approach can be made with more data at smaller angles (e.g., the s dependence at $\vartheta_{c.m.}=60°$) and with study of the other elastic reactions such as $\pi p \rightarrow \pi p$. In this last case, there is a hint of oscillatory s dependence in existing data which would be interesting to compare with nuclear targets.

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[13]We take $A=12$, 27, 64, and 207 for comparison with C, Al, Cu, and Pb. The comparison with Li, for which two data points have been measured, is also consistent but adds little because the limit $A \gg 1$ does not clearly apply.

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