CERN-TH. 4476/86

SWELLING NUCLEONS IN QCD

J.V. Noble^{*)}

CERN - Geneva
and
University of Geneva, Charlottesville VA 22901

A B S T R A C T

The analogy between the QCD vacuum - supposed to be a colour-magnetic superconductor - and the superconducting state, is exploited to estimate the reduction in average string tension experienced by hadrons in nuclei. This reduction leads to larger radii and softer one-body momentum distributions of quarks in nuclei, in qualitative agreement with experiment. The mechanism avoids colour van der Waals forces that would otherwise arise from direct gluonic interactions between nucleons.

^{*)} This work has been supported in part by the U.S. National Science Foundation.

Recent data on both inclusive and exclusive quasielastic electron scattering [1, 2] are best interpreted by assuming the nucleon's structure is renormalized by dense nuclear matter [3]. Subsequent data on deep-inelastic scattering of electrons and muons from heavy nuclei [4, 5, 6] confirm this interpretation at the level of quark wavefunctions [7]. It is now generally agreed that a meson + baryon phenomenological field theory can be an appropriate low-energy limit of QCD. Such theories generally predict [3] renormalizations like those observed experimentally. However, it seems cumbersome to have to describe the softening of the single-quark momentum distribution (of a nucleon embedded in a nucleus, relative to the free nucleon's) by first going to a meson description and then translating back to quarks. A method for arriving at this result directly would seem preferable.

The object of this note is to suggest how, within the framework of the usual view of confined QCD, nearby nucleons can modify others without generating color van der Waals forces. That is, we need to assume neither gluon nor quark interchange between color singlets.

We begin with a standard view of the QCD vacuum as a condensate of gluons and qq pairs [8]. We do not pretend to know how to solve the formidable problem of producing a convincing model of the QCD condensed vacuum (from here on referred to as "the VACUUM"). We merely suppose the VACUUM is a superconductor of color magnetic field and expels the color electric field [9]. Only if enough energy per unit volume is supplied to break down the VACUUM to the "normal" or decondensed state (called "the vacuum" henceforth) can color electric flux propagate. Thus, the region between classical color charges (e.g. quark and antiquark) is a finite-volume region of vacuum embedded in VACUUM, and the color electric field must be confined to this finite region between the sources (Fig. 1). As the separation between sources increases, the energy in the field plus vacuum grows (asymptotically it rises linearly) [10]. Naturally this picture requires that isolated systems have zero net color charge, otherwise their energy would be infinite.

Consider well-separated (so we can ignore end effects) classical charge g and anticharge, $-g$. Their total energy will be

$$U = \int_{-\infty}^{\infty} dz \int d^2b \left[\frac{1}{2} \mathbf{E} \cdot \mathbf{E} + C + j^0 \phi \right] \quad (1)$$

where j^0 is the color charge density, ϕ is the color electric potential defined by (we neglect the non-Abelian part)

$$E = - \nabla \phi, \quad (2)$$

and C is the energy per unit volume of vacuum relative to VACUUM. Suppose we have sources $\pm g\delta(r \pm \frac{1}{2}a)$: then at short distance $|a|$ between the sources we expect the interaction energy to be dominated by the coulombic term

$$V_{\text{coul}} = -\frac{g^2}{4\pi|a|}, \quad (3)$$

and at long distances we expect the interaction energy to rise like $\sigma|a|$, where σ is the string tension. From studies of mesons composed of heavy quarks, σ is empirically 1.2 GeV/fm. The asymptotic string tension can be calculated by noting that if $|a|$ is large, the field E within the tube becomes $E(b)z$; the energy minimum principle is then (ignoring end effects)

$$-\frac{\delta}{\delta E(b)} \int_0^\infty db \, b \left[\frac{1}{2} E^2(b) + C\theta(E^2(b)) \right] = 0 \quad (4)$$

where $E(b) \geq 0$ and

$$\int_0^\infty db \, b \, E(b) = -\frac{g}{2\pi}. \quad (5)$$

Clearly, E must be constant across a tube of finite radius β ,

$$E(b) = E_0 \theta(\beta - b) \quad (6)$$

and the string tension may be rewritten

$$\sigma(\beta) = -\frac{g^2}{2\pi\beta^2} + C \pi\beta^2. \quad (7)$$

We now minimize with respect to $x = \pi\beta^2$, to get

$$\sigma_{\min} = \left[8\pi C - \frac{g^2}{4\pi} \right]^{1/2} \quad (8)$$

We find that C must be of order 300 MeV/fm^3 ($g^2/4\pi \approx 1$) in order to obtain the correct asymptotic string tension. This seems reasonable, since the volume of light-quark hadrons must be of order $2\text{--}4 \text{ fm}^3$, and such a value of C would produce the correct mass scale (keeping in mind the attractive part, Eq. (3)). Note that the tube radius turns out to be $\beta \approx 0.7 \text{ fm}$, which is not much like a string (this remark is not original [8, 10]).

Now, supposing that the color electric flux within nucleons behaves qualitatively like the classical picture described above, let us imagine what happens when two identical nucleons come near each other (Fig. 2).

- i. The quarks will find their possible phase space restricted by the Pauli principle. For only two nucleons, this is not much of a restriction (because of the three-fold color degree of freedom) and does not raise the energy substantially [11].

- ii. Although $SU(3)_C$ gauge field theory leads to cubic and

quartic couplings, suppose the energies associated with such nonlinearities are also not large, so that superposition holds approximately. That is, we may suppose the color electric fields of one nucleon interpenetrate those of another, with little energy penalty. In bag model language, the bags can overlap, more-or-less freely. (This is not the same as imagining multi-quark bags, however, since it is by no means obvious that a rearrangement into a system with -- say -- 6 quarks in an $[8 \times 8]_0$ state will

have lower energy than two interpenetrating 3-quark systems in the $[0 \times 0]_0$ state.)

When the color electric fields of a hadron interpenetrate those of another, they pass -- in part -- through vacuum, rather than VACUUM. Suppose the overlap volume is $\delta\Omega$: then the energy of vacuum relative to VACUUM must be reduced by, roughly,

$$C \rightarrow C - C\delta\Omega/2V_N, \quad (9)$$

to avoid double-counting. Here V_N is the volume occupied by the color electric fields binding the nucleon.

What is the overlap volume $\delta\Omega$? In a nucleus the volume per nucleon -- or "specific volume" -- is

$$V_0 = 4\pi r_0^3/3$$

where r_0 is the usual nuclear radius parameter $R = r_0 A^{1/3}$, about 1.1 to 1.2 fm. That is, r_0 is the Wigner-Seitz radius. If the density is not too high -- so we can neglect 3- and 4-body overlaps -- the vacant fraction of volume in a specific volume V_0 occupied by a nucleon is

$$f = 1 - V_N/V_0 ;$$

hence the occupied fraction is $1 - f$. Then the overlap volume for two nucleons is -- on the average --

$$\delta\Omega = V_N(1 - f) = V_N^2/V_0 .$$

The fractional reduction in vacuum-VACUUM energy C is then

$$-\frac{\delta\Omega}{2V_N} \approx -\frac{1}{2} (r_N/r_0)^3 \approx 0.15 - 1.0 , \quad (10)$$

where we have taken the nucleon "glue" radius from 0.8 fm to 1.4 fm ($= 2 \times 0.7$ fm), and the Wigner-Seitz radius r_0 from 1.1 to 1.2 fm. The lowering of the effective vacuum energy C will lower the energy of the whole system since C is the only dimensional parameter. Except for the attractive coulombic term (3) (and perhaps attractive nonlinear interactions) it would be advantageous for the radii of nucleons in nuclei to expand as much as

possible, and thereby overlap as much as possible.

The upshot is that the effective string tension of a nucleon in a nucleus is reduced from that of an isolated nucleon. The relative decrease in string tension implied by Eq. (10) is 7% to 30%, implying a similar decrease in the nucleon mass and corresponding increase in its radius. (The radius could increase by 40%.)

We emphasize that because we have linearized a fundamentally nonlinear problem, the estimates must be considered rough. A better calculation taking into account end effects, the more complex color electric field of the 3 quark system, and the nonlinear self interactions of gluons -- even within the framework of the particular vacuum structure postulated here -- needs to be done and is being planned. Nevertheless the estimate (10) agrees encouragingly, both in direction and sign, with the general range of the renormalization effects seen in quasielastic and deep-inelastic scattering from nuclei.

There appears to be some feeling that the EMC effect [4] is essentially trivial, since it "merely" reflects an increase in the confinement radius of valence quarks -- known already from intermediate energy quasielastic electron scattering -- through the corresponding decrease in their average momentum. This is far from the case: it has proven surprisingly difficult -- as the many papers on the subject will attest -- even to pin down the phenomenology of the EMC effect and identify it unequivocally with a change in the valence quark momentum distribution of the nucleon [7]. Having made such an identification, it was clear that tracing the change in structure (owing to the nuclear environment) could shed light on the confinement mechanism. This paper has shown that with some not too outrageous assumptions, a popular model of confinement leads to an effect in the right direction and of the right order. The EMC effect can now be considered a crucial test that other proposed confinement mechanisms must pass to be credible.

I am grateful to CERN for its hospitality during the course of this work, and to V. Balunyi and G. Kane for many helpful discussions.

References

1. R. Altemus, *et al.*, Phys. Rev. Lett. **44** (1980) 965.
M. Deady, *et al.*, Phys. Rev. **C28** (1983) 631.
Z. Meziani, *et al.*, Phys. Rev. Lett. **52** (1984) 2180;
54 (1985) 1233.
2. J. Morgenstern, Nucl. Phys. **A446** (1985) 315c.
3. J.V. Noble, Phys. Rev. Lett. **46** (1981) 412;
Nucl. Phys. **A329** (1979) 354.
4. J.J. Aubert, *et al.* (EMC), Phys. Lett. **123B** (1983) 275.
5. G. Bari, *et al.* (BCDMS), Phys. Lett. **163B** (1985) 282.
6. A. Bodek, *et al.* (SLAC), Phys. Rev. Lett. **52** (1984) 727.
7. J.V. Noble, CERN preprint CERN-TH.4441/86, and submitted to Nucl. Phys. **B**.
8. P. Hasenfratz and J. Kuti, Phys. Rep. **40C** (1978) 75. (These authors give an early discussion of classical flux tube models like those leading to Eq. (4)-(8).)
L.J. Reinders, H. Rubinstein and S. Yazaki, Phys. Rep. **127C** (1985) 1.
S. Weinberg, in A Festschrift for I.I. Rabi, ed. L. Motz (New York Academy of Sciences, New York, 1977) p. 185.
9. S. Mandelstam, Phys. Rep. **23C** (1976) 245.
C.G. Callen, R. F. Dashen and D.J. Gross, Phys. Rev. **D20** (1979) 3279.
10. S.L. Adler and T. Piran, Rev. Mod. Phys. **56** (1984) 1.
11. C. DeTar, Phys. Rev. **D17** (1978) 323.
See also M. Harvey, in Short-Distance Phenomena in Nuclear Physics, ed. D. Boal and R. Woloshyn (Plenum Press, New York, 1983) p. 101.

Figure Captions

1. A flux tube picture of the region between a quark and antiquark.
2. Flux-tube picture of two overlapping hadrons, with the included volume $\delta\Omega$ shaded.

Fig. 1

Fig. 2