MANY-NUCLEON FORCES AND CURRENTS DERIVED FROM QUARK DEGREES OF FREEDOM+

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1. Introduction

Nuclear physics is still on the search for a workable and quantitatively successful microscopic theory of nuclear phenomena. This paper wants to contribute to this on-going search.

The traditional concept of nuclear theory views the nucleus as a system of elementary nucleons interacting through instantaneous two-nucleon forces according to the rules of nonrelativistic quantum mechanics. This many-body problem of the nucleus can be solved for three nucleons with satisfying numerical accuracy. Its solution demonstrates: The classic theory of nuclear phenomena is qualitatively successful for bound and scattering states, but it does not yield[1)] quantitative agreement with experimental data. Its successes warn to give up the traditional concept as wrong altogether, but the persisting discrepancies with experiment necessitate corrections.

The prominent problem of traditional microscopic nuclear structure is the theoretical description of binding energies and electromagnetic (e.m.) properties. The standard corrections take into account the virtual excitation of a nucleon to an isobar and mesonic exchange currents. These special correction mechanisms yield[2)] encouraging improvements of the theoretical results, but - in the three-nucleon system, the test case for the more complicated heavier nuclei - they are unable to fully cure the existing discrepancies. However, the nucleon is itself a composite system consisting of quarks, antiquarks and gluons. This substructure of the nucleon only plays an insignificant role for the correction mechanisms considered so far. Because of the remaining discrepancies it therefore appears reasonable to investigate corrections of traditional microscopic nuclear structure, which only arise in many-nucleon systems and which are due to the internal nucleonic degrees of freedom. This is the motivation for the work reported on in this paper.

In a quantum-chromodynamical (QCD) description most nuclear problems are phenomena of strong quark-gluon coupling. The absence of an explanation, why nature chooses colorless quark clusters only, i.e., why quarks are confined, and the absence of a technically simple and micro-

+Talk presented by P.U. Sauer

scopically convincing phenomenology of confinement appear to be the stumbling blocks for a successful treatment of the internal nucleonic degrees of freedom in nuclear physics. The employed methods for generating confinement by "fiat" are bag models with relativistic, practically massless quarks and potential models of nonrelativistic quarks with an effective constituent mass. In the latter models gluon degrees of freedom are frozen into quark-quark potentials and only the quark degrees of freedom are kept.

The internal nucleonic quark degrees of freedom are usually envoked to understand[3)] known nucleonic properties, e.g., e.m. form factors, meson-nucleon couplings and the nucleon-nucleon interaction, on a more fundamental level. However, these nucleonic properties are only the input parameters for the nuclear many-body problem, which then remains to be solved in a Hilbert space of nucleons only as in traditional microscopic nuclear structure. In this way qualitatively new results for the many-nucleon systems and a cure of existing discrepancies between theory and experiment are not to be expected: The dependence of nuclear-structure results on a variation of these input parameters has already been tested phenomenologically, it is the dependence on the off-shell properties[4)] of the two-nucleon potential.

The present work is complementary to the procedure just mentioned. It aims at those many-nucleon phenomena which can only arise from the quark degrees of freedom in nucleons. The phenomena to be investigated are based on nonnucleonic hidden-color six-quark states: When the three-quark clusters of two nucleons in the nucleus have a nonvanishing overlap, six-quark configurations arise through gluon exchange and the indistinguishability of identical quarks which can no longer be interpreted as states of two physically separable baryons. For these nonnucleonic six-quark states the quark property color is essential. Whereas all physically observed quark systems are colorless, the colorless six-quark states may also contain spatially localized substructures with nonvanishing color[5)], hidden-color six-quark states. Such quark states have been employed in the two-nucleon system for the description of elastic nucleon-nucleon scattering[6)] at low energies and in the analysis of the e.m. deuteron properties[7)]. Such quark states yield the basis for the attempted quark interpretation of the structures seen in nucleon-nucleon scattering at intermediate energies as dibaryon resonances[8)]. The present work extends their use to many-nucleon systems. In the nuclear medium these six-quark states yield new effects: (i) Many-nucleon forces arise when a hidden-color six-quark state interacts with the colorless three-quark clusters of other nucleons. (ii) The hidden-color six-quark states have a spin-isospin structure, which differs from that of two interacting nucleons, and therefore also a different coupling to the photon field. Thus, a new type of many-nucleon currents arises. This paper formulates these new quark phenomena in nuclear physics, i.e., many-nucleon forces and currents, as corrections for traditional microscopic nuclear structure and tests their consequences for understanding nuclear properties by the example of the bound three-nucleon system.

2. Two-nucleon interaction model with hidden-color six-quark states

The aim of this section is to develop a model for the two-nucleon interaction with novel characteristics. The description of the interaction should remain traditional and macroscopic, i.e., it should use the coordinate $\vec{r}$, i.e., the separation between two three-quark clusters, the total spin S and isospin I as dynamic variables. However, the description should also contain physically relevant features of the composite nature of nucleons, i.e., of subbaryonic degrees of freedom. The latter goal will be acchieved by coupling channels with nonbaryonic quark content to two-baryon channels traditionally employed.

We view two baryons in interaction as a system of six quarks interacting nonrelativistically through instantaneous quark-quark potentials. The interaction between two quarks is parametrised in the usual form $\vec{\lambda}_i \cdot \vec{\lambda}_j\ V(i,j)$. The $\vec{\lambda}_i$ are the eight component vectors of SU(3) generators in color space and the subscripts i and j denote the particle labels of the quarks. The potential V(i,j) contains the Fermi-Breit one-gluon exchange[9], nonrelativistically approximated and reduced to its local central pieces and a harmonic confinement potential, i.e.,

$$V(i,j) = \frac{1}{4}\hbar c \alpha_s \left[\frac{1}{|\vec{r}_i - \vec{r}_j|} - \pi \left(\frac{\hbar}{mc}\right)^2 \delta(\vec{r}_i - \vec{r}_j)\left(1 + \frac{2}{3}\, \vec{\sigma}_i \cdot \vec{\sigma}_j\right) \right] - \frac{3}{16} k (\vec{r}_i - \vec{r}_j)^2. \quad (2.1)$$

The emerging quark hamiltonian of internal motion has the form

$$H = \sum_i \left(\frac{p_i^2}{2m} + mc^2\right) - \frac{\left(\sum_i \vec{p}_i\right)^2}{2 \sum_i m} + \sum_{i<j} \vec{\lambda}_i \cdot \vec{\lambda}_j\ V(i,j). \quad (2.2)$$

It only makes sense in a restricted space of states with vanishing total color.

The hamiltonian (2.2) is first applied to three-quark systems in order to determine its parameters. The three-quark systems are assumed to consist of up and down quarks only and to cluster around a center in space located, say, at $\vec{R}_a$. The three-quark wave function of the baryon A

$$<\vec{r}_1\vec{r}_2\vec{r}_3|A_q\vec{R}_a> = <\vec{r}_1|\vec{R}_a><r_2|\vec{R}_a><\vec{r}_3|\vec{R}_q>|A_q> \quad (2.3)$$

has a spatially symmetric part, the normalized product of harmonic oscillator wave functions $<\vec{r}_1|\vec{R}_a><\vec{r}_2|\vec{R}_a><\vec{r}_3|\vec{R}_a>$ in the lowest orbit with the oscillator parameter α, and a totally antisymmetric spin-flavor-color part $|A_q>$. Given the spatial wave function there are six different types of three-quark clusters. Two of them are color singlets, they correspond to the nucleon (N) and the delta-isobar (Δ). Three of them are color octets (0,0' and 0"), which differ by isospin and spin, i.e., $(IS)=(\frac{1}{2}\ \frac{1}{2}),(\frac{1}{2}\ \frac{3}{2}),(\frac{3}{2}\ \frac{1}{2})$ respectively. The sixth is a color decuplet. The three-quark cluster wave functions are given in ref. (10). The parameters of the quark hamiltonian (2.2), i.e., the quark mass m, the strong coupling constant α_s, the confinement parameter k and the oscillator parameter α for the spatial wave functions are fixed by properties of the physical baryons nucleon and delta using the nucleonic magnetic moments and the nucleon-delta mass difference $M_\Delta c^2 - M_N c^2$. The masses of the nucleon and the Δ-isobar are required to be stationary under variation of the oscillator size parameter α. The colored three-quark clusters do not have a realization as isolated physical objects. However, the three color octets can serve as building blocks for obtaining all colorless six-quark configurations. This is their function in the present context. The color-decuplet cluster does not form a colorless six-quark state with any other three-quark cluster due to SU(3) coupling rules and therefore will not be discussed any further.

Next the quark hamiltonian (2.2) is applied to six-quark systems. We are very well aware of the fact that this hamiltonian cannot describe the intermediate- and long-range attraction of the two-nucleon potential. In fact it yields a wrong long-range van-der-Waals force[11)], which is contradicted by the low-energy nucleon-nucleon scattering data. The hamiltonian has proven to be successful for short-range phenomena, for which confinement does not play a crucial role. We therefore use it for a description of the dynamic evolution of the six-quark system at small cluster separations and shall phenomenologically correct its faulty aspects at large distances. For the description we assume that three quarks always remain clustered around their individual centers. Admittedly, this is an asumption hard to justify.

We describe the six-quark system with the tools of the generator coordinate method (GCM)[12)]. The physically relevant generator coordinate is the distance $\vec{r}$ between two centers $\vec{R}_a$ and $\vec{R}_b$, i.e., $\vec{r} = \vec{R}_a-\vec{R}_b$, around which the two three-quark clusters A_q and B_q are built. The used basis states are[10)]

$$|\beta\vec{r}> = |A_q\ B_q\ ISC\ \vec{r}> \qquad (2.4a)$$

$$|A_q B_q ISC\vec{r}> =$$

$$\frac{1}{\sqrt{2}}\bar{\mathcal{A}}\sum_{\Theta}\mathcal{N}(\Theta)|\Theta\vec{r}>\frac{1}{\sqrt{2}}\left[|A_q B_q ISC>+(-1)^{I-I_A-I_B+S-S_A-S_B+\Theta}|B_q A_q ISC>\right] \qquad (2.4b)$$

The basis states are colorless, i.e., C = 1, and totally antisymmetrised. They are constructed from orbital states $|\Theta\vec{r}>$ of definite permutation symmetry Θ, Θ being 3, 2, 1 and 0 for the Young tableaux [6], [51], [42] and [33] respectively, and from spin-flavor-color states $|A_qB_qISC>$ in which the individual cluster states $|A_q>$ and $|B_q>$ are coupled to total isospin I, total spin S and total color C. Antisymmetrization within the clusters and cluster exchange is carried out beforehand, thus $\bar{\mathcal{A}} = \frac{1}{10}\left[I-P_{14}-P_{15}-P_{16}-P_{24}-P_{25}-P_{26}-P_{34}-P_{35}-P_{36}\right]$ denotes the remaining antisymmetrization between clusters only. The channel label β stands for the cluster content A_q and B_q and the total quantum numbers; isospin and spin projections are suppressed. There are two-baryon channels β_B, channels $\{N_qN_qIS\}$ with purely nucleonic content and channels $\{N_q\Delta_qIS\}$, $\{\Delta_qN_qIS\}$ and $\{\Delta_q\Delta_qIS\}$, well-known from usual corrections of the classic picture of microscopic nuclear structure, in which one or both nucleons are turned into a Δ-isobar. However, there are also channels β_C with color-octet content, i.e., $\{0_q0_qIS\}$, $\{0_q0_q'\,IS\}$, $\{0_q0_q''\,IS\}$, $\{0_q'\,0_qIS\}$, $\{0_q'\,0_q'\,IS\}$, $\{0_q'\,0_q''\,IS\}$, $\{0_q''\,0_qIS\}$, $\{0_q''\,0_q'\,IS\}$, $\{0_q''\,0_q''\,IS\}$, which appear as novel channels of hidden color. They arise from the quark description and are no realization of simple two-baryon states. The still unspecified weights $\mathcal{N}(\Theta)$ and phases $(-1)^{\Theta}$ are given in ref. (10). The basis $\{|\beta\vec{r}>\}$ is nonorthogonal and noncomplete in the six-quark Hilbert space.

The GCM ansatz for the dynamic many-quark wave function $|\phi_E>$ of the collective motion in the degree of freedom associated with the generator coordinate $\vec{r}$ is a linear superposition of the noncomplete basis states $\{|\beta\vec{r}>\}$, i.e.,

$$|\bar{\Phi}_E> = \sum_{\beta}\int d^3r\,|\beta\vec{r}><\beta\vec{r}|\phi_E> \qquad (2.5)$$

Diagonalizing the hamiltonian H of equ. (2.2) in the subspace spanned by the basis states leads to the Hill-Wheeler integral equations for the weight function $<\beta\vec{r}|\phi_E>$ which contains the information on the dynamics of collective motion

$$\sum_{\beta'} \int d^3r' [<\beta\vec{r}|H|\beta'\vec{r}'> - E <\beta\vec{r}|\beta'\vec{r}'>] <\beta'\vec{r}'|\phi_E> = 0 \qquad (2.6a)$$

$$[\bar{H} - E\bar{N}] \; |\bar{\phi}_E> = 0 \qquad (2.6b)$$

Equ. (2.6b) is the operator abbreviation of the Hill-Wheeler equation in the subspace $\{\beta\vec{r}>\}$ of the Hilbert space. Let $\bar{X}$ be a - not necessarily hermitian and certainly non-unique - square-root of the overlap operator $\bar{N}$ in the considered subspace, i.e., $\bar{X}^+\bar{X} = \bar{N}$, then the Hill-Wheeler equation can be formally rewritten as

$$\sum_{\beta'} \int d^3r' <\beta\vec{r}|\mathcal{H}|\beta'\vec{r}'><\beta'\vec{r}'|\Psi_E> = E <\beta'\vec{r}'|\Psi_E> \qquad (2.7)$$

with $\mathcal{H} = (\bar{X}^{-1})^+ \bar{H} \bar{X}^{-1}$ and $|\Psi_E> = \bar{X}|\bar{\phi}_E>$. This is a conventional Schrödinger equation in the collective variable $\vec{r}$ which couples channels of different cluster content β. The states satisfy the usual orthonormality condition $<\Psi_E'|\Psi_E> = \delta(E'-E)$. They can be regarded as a traditonal two-baryon wave function with a direct probability interpretation on which boundary conditions for the scattering and bound-state situations can be imposed as usual. Physicswise, the Schrödinger equation (2.7) describes an interaction model which couples two-baryon channels β_B of separable colorless three-quark clusters with hidden-color channels β_C. The latter channels are spatially confined. Thus, the coupling between two-baryon channels and channels of hidden-color and the propagation in the latter channels occur at small relative distances $\vec{r}$.

The nonrelativistic one-gluon exchange hamiltonian of equ. (2.2) has been successful in describing small-distance quark dynamics. Thus the coupling between baryonic and nonbaryonic six-quark channels $<\beta_B\vec{r}|\mathcal{H}|\beta_C\vec{r}'>$ and the propagation in the nonbaryonic channels $<\beta_C\vec{r}|\mathcal{H}|\beta_C'\vec{r}'>$, derived from the hamiltonian (2.2) should be physically reliable. In contrast, the hamiltonian (2.2) cannot account for the intermediate and long range of the two-baryon interaction. We therefore have to repair this defect and replace the part $<\beta_B\vec{r}|\mathcal{H}|\beta_B'\vec{r}'>$ of the derived effective hamiltonian by a traditional one-boson exchange description, i.e.,

$$\langle\beta_B\vec{r}|\mathcal{H}|\beta_{B'}\vec{r}'\rangle = -\frac{\hbar^2}{2\mu_B}\Delta_{\vec{r}}\,\delta(\vec{r}-\vec{r}')\,\delta_{BB'} + \langle\beta_B\vec{r}|V_{OBE}|\beta_{B'}\vec{r}'\rangle \tag{2.8}$$

In equ. (2.8) μ_B is the reduced mass of the channel and V_{OBE} the one-boson exchange interaction between nucleons and isobars.

Thus, a two-nucleon force model is obtained in equs. (2.7) and (2.8), which remains in the traditional framework by using baryon separation, spin and isospin as dynamic variables, as well as an ordinary Schrödinger equation for the description of its time evolution. The force model yields an interaction between nucleons and Δ-isobars, but it is amended through coupling to nonbaryonic channels of hidden color for subnucleonic degrees of freedom. By the phenomenological replacement (2.8) the force model becomes a truely realistic one.

3. The interaction with hidden color channels

In this section two nucleons whose dynamics are governed by the force model of equs. (2.7) and (2.8) are emersed in a many-nucleon system and are subjected to external fields. The interaction with other nucleons and with the external sources is discussed. Novel many-nucleon forces and currents arise. The many-nucleon system is described in a Hilbert space of A baryons with a sector in which two baryons are turned into a nonbaryonic bound-state channel of hidden-color characteristics β_C. The baryons may be either nucleons or Δ-isobars. The considered Hilbert space is diagramatically described in fig. 1.

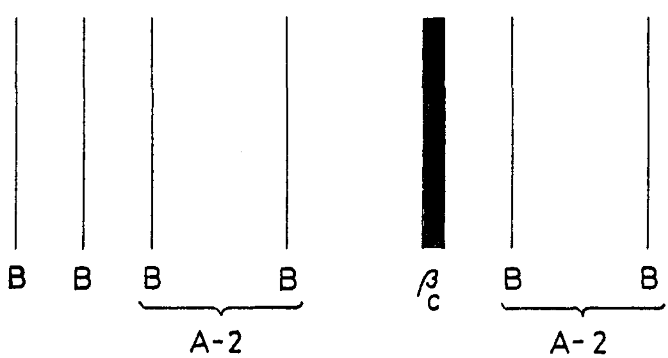


Fig. 1: The Hilbert space of a nucleus with baryon number A. It has two sectors. The first one contains A ordinary baryons B, in the second sector two baryons are excited to a hidden-color state of characteristics β_C.

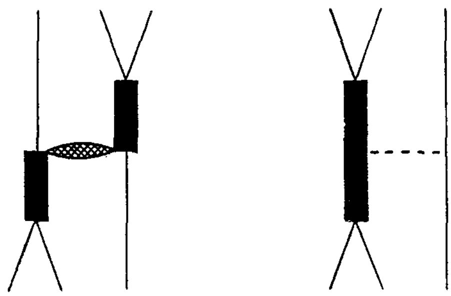


Fig. 2: Three-baryon forces. In the first type (a) the three-baryon force is due to color transport from one pair to another. It is of short range. The second type (b) is of long range arising from meson exchange between a hidden-color pair and an ordinary baryon.

Firstly the quark hamiltonian (2.2) can switch the hidden-color structure β_C from one pair of interacting baryons to another. This process yields a three-nucleon force, in general a many-baryon force, as diagramatically shown in fig. 2a. It can occur only when all three nucleons are close to each other.

When the quark system is subjected to external sources, the quark hamiltonian H of (2.2) has to be extended for additional couplings, i.e.,

$$H_{\phi A} = H + \sum_i g \vec{\sigma}_i \cdot \vec{\nabla}_x \vec{\tau}_i \cdot \vec{\phi}(\vec{x}) + \sum_i \frac{1}{c} j_{i\mu}(\vec{x}) A^{\mu}(\vec{x}) \tag{3.1}$$

Clearly, quarks are the carriers of the e.m. current of baryons and the third term is the appropriate interaction with an external e.m. source of four-vector potential $A^{\mu}(\vec{x})$. Term two describes the coupling of quarks with a mesonic field. In equ. (3.1) the mesonic field is treated as elementary. This approximation neglects the internal quark-antiquark structure of mesons. As an example the isovector pion field $\vec{\phi}(\vec{x})$ is used in (3.1). It is assumed to originate from those nucleons of the many-nucleon system not described by their quark degrees of freedom through the hamiltonian (2.2) and (3.1). Due to term two of $H_{\phi A}$ the six-quark bound-state structures of hidden-color characteristics β_C, when emersed in a many-nucleon system, can also interact with the colorless three-quark clusters of other baryons in the traditional way. The interaction is of one-boson exchange type. Compared to the first process of fig. 2a the meson coupling in $H_{\phi A}$ yields a second three-nucleon force diagramatically described in fig. 2b which is of longer range and therefore might be physically more important.

Thirdly, the six quark bound-state channels β_C interact with photons. Compared to two nucleons in interaction the spatial location and the spin - isospin structures of the hidden-color channels are different. Thus, they yield e.m. two-nucleon, generally many-baryon, currents of novel characteristics abbreviated by the diagram of fig. 3.

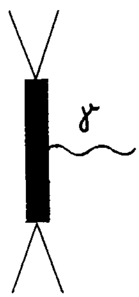


Fig. 3: Two-baryon exchange current arising from the coupling of the photon to a hidden-color pair.

The basic result of this section is: When the two-nucleon system is described by its quark degrees of freedom, corrections for the traditional picture of the nucleus arise. The interaction of the two nucleons with other nucleons in the many-nucleon system yields novel many-nucleon forces, its interaction with an external e.m. probe yields novel many-nucleon exchange currents. The remainder of this paper discusses the practical applicability of these many-nucleon forces and currents in nuclear structure and estimates their effect for the three-nucleon bound states.

4. Simulation of hidden-color six-quark channels by two confined color-excited baryons

The theoretical frame work of sections 2 and 3 is practicable and could be employed, e.g., in the three-nucleon bound state, but, admittedly, it is technically demanding. For first estimates of the importance of the resulting many-nucleon forces and currents we used the following simplification for the hidden-color six-quark states β_C. They are assumed to be approximated by two point baryons C which carry the new quantum number color besides the quantum numbers isospin and spin, i.e., C being O, O' or O" of section 2. The color-octet baryon spans the eight-dimensional representation of color SU(3). Each color-octet baryon requires a color-octet partner in the many-nucleon system to whom it is confined and with whom together it forms a colorless bound two-baryon state. Thus the Hilbert space which describes a physical A-nucleon system consists - besides the ordinary part of colorless baryons - of a sector with two color-octet baryons coupled to a color-singlet. In this approximation the A-nucleon system remains in its microscopic description an A-baryon system. The quark degrees of freedom are entirely eliminated, but they have given rise to baryon states with new properties. The Hilbert space is described in fig. 4. It is to be interpreted as an approximative version of the Hilbert space of fig. 1.

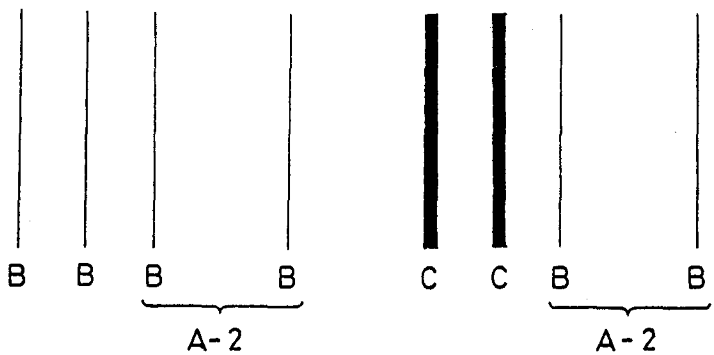


Fig. 4: The Hilbert space with the hidden-color pair simulated by two color-octet baryons C.

In the Hilbert space with color-octet baryons the dynamics of the two-baryon system is approximately described by the following coupled-channel Schrödinger equation.

$$\left[(M_B(1)+M_B(2)-2M_N)c^2 - \frac{\hbar^2}{2}\left(\frac{1}{M_B(1)} + \frac{1}{M_B(2)}\right)\Delta_{\vec{r}}\right]<\beta\vec{r}|\Psi_E> +$$

$$+ \sum_{\beta'} \int d^3r'<\beta\vec{r}|V|\beta'\vec{r}'><\beta'\vec{r}'|\Psi_E> = E<\beta\vec{r}|\Psi_E>. \tag{4.1}$$

In the channels with nucleons and isobars the interaction used is the one of equ. (2.8). The coupling to color-excited baryonic configurations and the potentials in color-excited channels are identified with the generator-coordinate matrix elements of equ. (2.7). These potentials are only effective at small relative distances between baryons. They are used in a local and - neglecting the necessary transformation $\overline{X}$ to a unit overlap kernel - in an adiabatic approximation.

$$\langle BB\ IS\ \vec{r}|V|CC\ ISF'\rangle = V^{IS}_{BB\to CC}(\vec{r})\ \delta(\vec{r}-\vec{r}\,') \tag{4.2}$$

$$V^{IS}_{BB\to CC}(\vec{r}) = \langle\{B_q B_q IS\}\vec{r}|H|\{C_q C_q IS\}\vec{r}\rangle \tag{4.2a}$$

$$\langle CC\ IS\ \vec{r}|V|CC\ IS\ \vec{r}\,'\rangle = V^{IS}_{CC\to CC}(\vec{r})\ \delta(\vec{r}-\vec{r}\,') \tag{4.3}$$

$$V^{IS}_{CC\to CC}(\vec{r}) = \langle\{C_q C_q IS\}\vec{r}|H|\{C_q C_q IS\}\vec{r}\rangle - 2\langle C_q|H|C_q\rangle \tag{4.3a}$$

The mass of the color-octet baryon needed for the formulation of the coupled-channel Schrödinger equation (4.1) is taken to be $M_C c^2 = \langle C_q|H|C_q\rangle$ The questionable notion of a mass of a color-octet baryon is not needed for a single baryon, but only for the propagation of a pair of color-octet baryons in the form $-\frac{\hbar^2}{2}\left(\frac{1}{M_C(1)} + \frac{1}{M_C(2)}\right)\Delta_{\vec{r}}$. This mass M_C drops out everywhere else. The identifications (4.2) and (4.3) specify the content of equ. (4.1) which describes the dynamics of the two-baryon system. Since the transition potentials (4.2) and (4.3) are localized to act at small relative distances, no long-range van-der-Waals force can arise. Admittedly, equ. (4.1) with these identifications is only a primitive approximation of the Born-Oppenheimer type for the full generator-coordinate formulation of equ. (2.7). Without solving (2.7) the validity of its approximative form (4.1) cannot be proven. We simply hope it to be sufficient for a first estimate of many-nucleon forces and currents arising from hidden-color configurations in nuclei.

The coupling potentials between the baryonic channels are abbreviated by the diagrams of figs. 5a to 5c. The Schrödinger equation (4.1) can be generalized in the usual way to many-nucleon systems. There the one-gluon exchange can carry color from one baryon to another. This process is indicated in fig. 5d. If mesons could be interpreted as elementary fields, color-octet baryons have a normal coupling to mesons and therefore a normal one-boson exchange interaction with color-singlet baryons as in fig. 5e. There is an ordinary coupling to the photon given in fig. 6. Quantitatively, the coupling of color-octet baryons with mesons and photons is determined by quark considerations in the same way as the couplings of nucleons and Δ-isobars with mesons and photons are related to each other. Through the couplings of figs. 5 and 6 many-nucleon

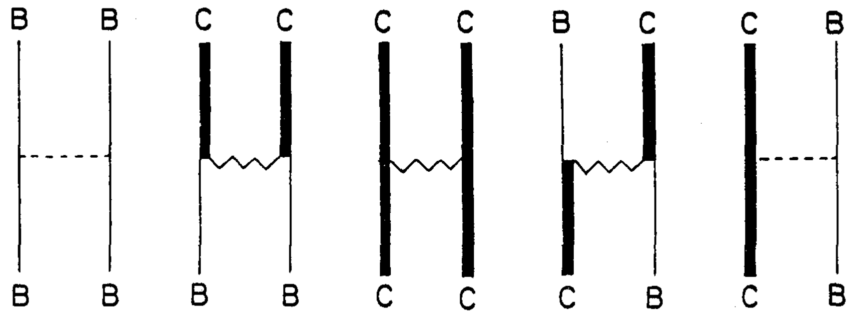


Fig. 5: Force model consisting of meson- and color-exchange between ordinary and color-octet baryons.

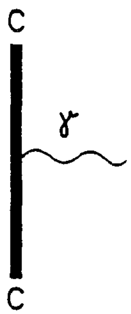


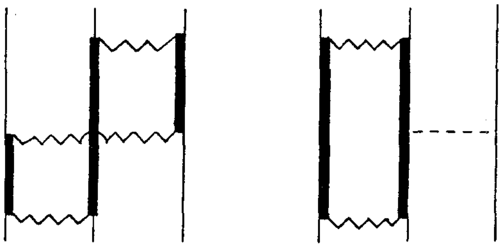
Fig. 6: Photon coupling to a color-octet baryon.

forces arise in the many-nucleon system as well as e.m. many-nucleon currents in the interaction with an e.m. probe. The three-nucleon forces are illustrated in fig. 7, they approximate the corresponding three-nucleon forces of fig. 2. The two-nucleon current is illustrated in fig. 8, it approximates the corresponding current of fig. 3.

5. Application in the three-nucleon bound state

The excitation of two nucleons to a hidden-color six-quark state at small relative distances yields novel many-baryon forces and currents. In this section their effect on the ^3H binding energy and the ^3He charge form factor is studied[10]. The hidden-color state is simulated as described in the previous section.

The hidden-color component in the wave function is calculated in perturbation theory. It is obtained in first order in the potential



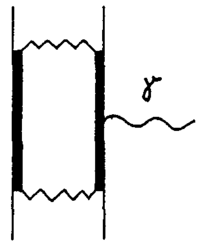


Fig. 7: The three-baryon forces of fig. 2 simulated with color-octet baryons.

Fig. 8: The two-baryon exchange current of fig. 3 simulated with color-octet baryons.

(4.2) from a purely nucleonic bound-state wave function. In the practical calculation the wave function of the Paris potential is used and together with the potential (4.2) a 0.5% probability for hidden-color components in the three-nucleon bound state arises.

Consistent with this rather small probability the effect of hidden-color configurations on binding energy and e.m. properties of the three-nucleon bound state is modest. In lowest nonvanishing order of perturbation theory the ^3H binding energy is affected by the three-baryon forces of fig. 7 and by the two-baryon dispersion diagramatically shown in fig. 9. The process of fig. 7a is of short range, since it is derived from short-ranged color exchange in each pair. The repulsion in the two-nucleon potential effectively screens this contribution to the three-nucleon force. Its effect is believed to be negligible and therefore has not been calculated. On the other hand the three-baryon force of fig. 7b contains a long-range meson exchange in one pair and should be much more important. Using π-, σ- and ρ-exchange as meson exchange the

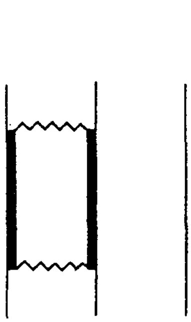


Fig. 9: Two-baryon dispersion weakening the three-baryon force of fig. 7b.

process of fig. 7b in fact yields an additional attraction of -0.8 MeV to the ^{3}H binding energy. This attraction is partly balanced by the two-baryon dispersion of fig. 9 resulting in repulsion of 0.1 MeV. According to this first primitive estimate the theoretical value of -7.4 MeV derived from the pure nucleonic Paris potential is increased by -0.7 MeV through hidden-color effects towards the experimental value of -8.5 MeV. In the same estimate the e.m. exchange current of fig. 8 yields a correction to the ^{3}He charge form factor. In fig. 10 the result including hidden-color contributions (solid curve) is compared to the pure nucleonic one of the Paris potential (dashed curve). The first minimum of the charge form factor is slightly shifted towards higher momentum transfers, whereas the slope around the 2. minimum is improved compared to the result for the pure nucleonic case.

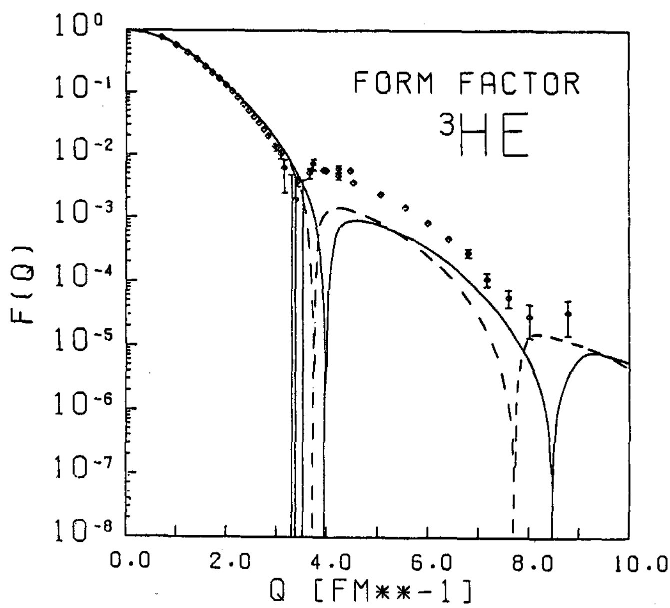


Fig. 10: Charge form factor of ^{3}He. The inclusion of hidden-color corrections (solid curve) is compared to the pure nucleonic result of the Paris potential (dashed curve).

6. Conclusion

The presented model is able to describe many-nucleon forces and currents derived from the quark substructure of nucleons more consistently and more quantitatively than it has been done before. The first results are clearly estimates only, all technical difficulties concerning the correct generator-coordinate treatment of the problem have been avoided in the present description. Many-baryon forces and currents appear small. We take this as a first indication that the traditional picture of the nucleus as an interacting system of nucleons, simple isobars and mesons is basically valid and requires only small corrections from subnuclear degrees of freedom.

References

1. "Mesons in Nuclei", ed. by M. Rho and D. Wilkinson (North-Holland, Amsterdam 1979); P.U. Sauer, Nuovo Cim. 76A (1983) 309
2. Ch. Hajduk, P.U. Sauer and W. Strueve, Nucl. Phys. A405 (1983) 581; W. Strueve, Ch. Hajduk and P.U. Sauer, Nucl. Phys. A405 (1983) 620
3. K. Johnson, Acta Physica Polonica B6 (1975) 2060; N. Isgur, "Soft QCD: Low Energy Hadron Physics with Chromodynamics" in "The New Aspects of Subnuclear Physics", ed. by A. Zichichi, Plenum Press 1980, p. 107
4. P.U. Sauer, "Phase-Equivalent Nonlocal Potentials" in "Nucleon-Nucleon Interactions", AIP Conference Proceedings 41, p. 195
5. M. Harvey, Nucl. Phys. A352 (1981) 301 and 326
6. A. Faessler, F. Fernandez, G. Lübeck and K. Shimizu, Phys. Lett. 112B (1982) 201
7. V.A. Matveev and P. Sorba, Lett. Nuovo Cim. 20 (1977) 435
8. A. Yokosawa, Phys. Rep. C64 (1980) 47; A. Th. M. Aerts, P.J.G. Mulders and J.J. de Swart, Phys. Rev. D17 (1978) 260
9. A. de Rujula, H. Georgi and S.L. Glashow, Phys. Rev. D12 (1975) 147
10. P.U. Sauer and U.-J. Wiese, to be published
11. M.B. Gavela, A. Le Yaouanc, L. Oliver, O. Pène, J.C. Raynal and S. Sood, Phys. Lett. 82B (1979) 431
12. H. Friedrich, Phys. Rep. C74 (1981) 209