IS THERE DIQUARK CLUSTERING IN THE NUCLEON?

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Abstract

It is shown that the instanton-induced interaction in \(qq\) pairs, iterated in \(t\)-channel, leads to a meson-exchange interactions between quarks. In this way one can achieve a simultaneous understanding of low-lying mesons, baryons and the nuclear force. The discussion is general and does not necessarily rely on the instanton-induced interaction. Any nonperturbative gluonic interaction between quarks, which is a source of the dynamical chiral symmetry breaking and explains the \(\pi - \rho\) mass splitting, will imply an effective meson exchange picture in baryons. Due to the (anti)screening there is a big difference between the initial ’t Hooft interaction and the effective meson-exchange interaction. It is demonstrated that the effective meson-exchange interaction, adjusted to the baryon spectrum, does not bind the scalar diquark and does not induce any significant quark-diquark clustering in the nucleon because of the nontrivial role played by the Pauli principle.

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I. INTRODUCTION

Speculations that instantons could induce diquark condensation in low temperature but high density quark matter \(^4\) have revived the interest in the diquark clustering in the nucleon. It is sometimes also argued that diquark condensation may occur even at moderate densities, for example in heavy nuclei. This problem is strongly related to the question of instanton induced diquark clustering in the nucleon. Indeed, the instanton-induced ’t Hooft interaction is strongly attractive for a quark-quark pair with quantum numbers \(T, J^P = 0, 0^+\) (scalar diquark). This raises expectations that it binds a scalar diquark and is responsible

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for the scalar diquark-quark structure of the nucleon [1,3]. This assumption is based on
the iteration of the ‘t Hooft interaction in the $qq$ $s$-channel. However, this picture of
the quark-quark interaction in baryons is only a small part of a more general one, based on the
effective meson-exchange interaction [4,5]. When the ‘t Hooft interaction is first iterated
in the $qq$ $t$-channel it inevitably leads to Goldstone boson exchange between constituent
quarks, which is drastically different from the initial (not iterated) ‘t Hooft interaction due
to the (anti)screening effects.

The latter effective meson-exchange interaction does not induce a bound scalar diquark,
nor an appreciable diquark-quark clustering in nucleon. This effective meson exchange
interaction is also the most attractive in $0,0^+$ $qq$ pairs, but the nature of this attraction is
very different from that of the ‘t Hooft interaction. This interaction, however, is not strong
enough to bind the scalar diquark. When it is combined with a confining interaction it binds
the diquark in the sense that there is no asymptotic state with two free constituent quarks,
though the mass of the scalar diquark is a few tens of MeV above the two-constituent-
quark threshold. There is no significant diquark clustering in the nucleon either, because
the nucleon is intrinsically a three-quark system and the fermionic-nature of the constituent
quarks plays an important role. If the subsystem of quarks 1 and 2 is in the $0,0^+$ state then
due to the antisymmetrization the quark pairs in the subsystems 1-3 and 2-3 are also partly
in the $0,0^+$ state. This implies that a strong attraction in $0,0^+$ quark pair contributes in all
quark subsystems simultaneously and makes the nucleon compact, but without appreciable
quark-diquark clustering.

This paper consists of two independent, but interrelated parts. In the first one we discuss
how the instanton-induced interaction (or some general nonperturbative gluonic interaction)
leads to the poles when it is iterated in the $qq$ $t$-channel. These pole contributions have an
evident meson-exchange interpretation. The latter meson-exchange interaction is drastically
different from the initial (bare) ‘t Hooft interaction which becomes strongly enhanced in the
channel of Goldstone boson exchange quantum numbers.

We also discuss the role of instantons in $\bar qq$ systems. There is no new wisdom in that
the nonperturbative gluonic configurations, e.g. instantons, induce the dynamical breaking
of chiral symmetry and explain the low-lying mesons. We include the latter discussion
only with the purpose of showing how the nonperturbative gluonic interaction both explains
mesons and at the same time leads to the effective meson exchange picture in the $qq$ sys-
tems. Through the latter it also explains the baryon spectra and the nuclear force. Our
discussion is rather general, and does not necessarily rely on the instanton-induced interac-
tion picture. Any nonperturbative gluonic interaction, which respects chiral symmetry and
induces the rearrangement of the vacuum (i.e. dynamical breaking of chiral symmetry), will
automatically explain the $\pi - \rho$ mass splitting and will imply a meson-exchange picture in
baryons.

The second part of this paper is devoted to a detailed study of diquark clustering in the
nucleon, based on the effective meson-exchange interactions in the baryons and the nucleon
wave functions obtained from the solution of the semirelativistic three-body Schrödinger
equation. We show that there is no appreciable diquark clustering in the nucleon and that
the effective meson-exchange interaction, which is adjusted to describe the baryon spectrum
[6], does not bind the scalar diquark nor the nucleon. However, when this interaction is
combined with the confining interaction, one finds a bound diquark but with a mass above
the two-quark threshold and very similar in magnitude to that obtained recently in lattice QCD \cite{7}. Nevertheless, as soon as the strength of the effective meson-exchange interaction is increased, not by a very big amount, it alone binds a nucleon, even without a confining force. While the contributions from the confining interaction to the nucleon mass are not small, the nucleon size, calculated with the confining interaction alone and in a full model that includes both confinement and effective meson exchange, is different. It is substantially smaller in the latter case, showing that there is indeed a soft interval between the scale when confinement becomes active, and the scale where chiral physics starts to work. However, for excited baryon states, which are much bigger in size, the role of confinement increases.

II. HOW ITERATION IN THE T-CHANNEL OF THE INSTANTON-INDUCED INTERACTION LEADS TO A MESON-EXCHANGE PICTURE AND (ANTI)SCREENS THE SHORT-RANGE BEHAVIOUR.

It has been shown in recent years that a successful explanation of light and strange baryon spectroscopy, especially the correct ordering of the lowest states with positive and negative parity, is achieved if the hyperfine interaction between constituent quarks $i$ and $j$ has a short-range behaviour which reads schematically \cite{4}:

$$-\vec{\lambda}_i^F \cdot \vec{\lambda}_j^F \vec{\sigma}_i \cdot \vec{\sigma}_j,$$

where $\lambda^F$ is a set of a flavor Gell-Mann matrices for $F = 1, ..., 8$ and $\lambda^0 = \sqrt{2/3}I$. This interaction is supplied by the short-range parts of Goldstone boson exchange (GBE)\cite{1}, vector-meson-like exchange and/or correlated two-pseudoscalar-meson-like exchange \cite{2}, etc.

It is sometimes stated that the instanton-induced ’t Hooft interaction in $qq$ pairs could also provide a good baryon spectrum as it contains a flavor- and spin-dependence and, iterated in the $qq$ s-channel, produces a deeply bound scalar diquark which makes the nucleon lighter than the $\Delta$ \cite{1,3}. A similar picture of a deeply bound scalar diquark has been advocated in a generalized Nambu and Jona-Lasinio (NJL) model \cite{5,6}. Then a baryon is constructed as an additive diquark-quark system or by solving “relativistic diquark-quark Faddeev equations” that take into account the quark exchange between the diquark and quark-spectator \cite{9,10}. In this section we show that such a picture of baryons, based on the iteration of the local 4-fermion interaction in the $qq$ s-channel is only a small part of a more general picture, based on the meson-exchange interaction. The reason is that when the ’t Hooft interaction (or generalized NJL one) is first iterated in the $qq$ t-channel, it inevitably leads to the effective meson-exchange between constituent quarks, which is drastically different from the initial (not iterated) 4-fermion local interaction due to (anti)screening effects. The difference is not only in the flavor- and spin-dependence, but sometimes also in the sign of the interaction.

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$^1$ $\pi$, $K$ and $\eta$ exchanges; due to the axial anomaly the $\eta'$ is not a Goldstone boson, but in the large $N_c$ limit it also becomes a Goldstone boson, and thus the coupling of $\eta'$ to a constituent quark should be essentially different from that of octet mesons.
To demonstrate this we use a simple $2 \times 2$ 't Hooft-determinant interaction for two light flavors (u and d), neglecting for the simplicity the tensor coupling term, which is suppressed by the factor $\frac{1}{4N_c} = \frac{1}{12}$ \[\text{cf.}\]. We also assume zero masses for the current u and d quarks in this section. For our illustrative purposes such an approximation is justified. This Hamiltonian reads:

$$H = -G[(\bar{\psi}\psi)^2 + (\bar{\psi}i\gamma_5\vec{\tau}\psi)^2 - (\bar{\psi}\vec{\tau}\psi)^2 - (\bar{\psi}i\gamma_5\psi)^2].$$ \hspace{1cm} (2)

The dimensional strength of the interaction $G$ as well as the ultraviolet cut-off scale $1/r_0$ can be related to parameters of the instanton liquid \[\text{cf.}\] (the dimensionless coupling constant is given by $G/4\pi r_0^2$). The interaction (2) is attractive in the scalar-isoscalar $\bar{q}q$ channel (the first term), leading to chiral symmetry breaking, or, which is related, to a massive $\sigma$-meson field and the constituent mass $m$ of quarks. This is readily obtained from the Schwinger-Dyson (gap) equation for a quark Green function in the Hartree-Fock approximation. The interaction in the $\bar{q}q$ pseudoscalar-isovector channel is driven by the second term of (2). It is so strong, that when it is iterated in the $\bar{q}q$ s-channel by solving the Bethe-Salpeter equation, see Fig. 1, it exactly compensates the $2m$-energy, supplied by the first term in (2), and thus there appear $T, J^P = 1, 0^-$ mesons with zero mass as deeply bound relativistic $\bar{q}q$ systems - Nambu-Goldstone bosons. The nonzero mass of the pseudoscalar mesons is brought about by the nonzero current quark mass as a perturbation, which is well illustrated by the current algebra results (Gell-Mann-Oakes-Renner relations). The first two terms in the Hamiltonian (3) form in fact the classical NJL Hamiltonian \[\text{cf.}\] and the statement above is a theorem, proved by Nambu and Jona-Lasinio many years ago. This scenario holds if the fixed strength of the interaction $G$ exceeds some critical level. In a more sophisticated derivation \[\text{cf.}\] the strength of the interaction $G$ is not fixed and should be determined after one gets the chirally broken phase.

The Hamiltonian (2) does not contain any interaction in $\bar{q}q$ pairs with vector meson quantum numbers. So, according to the scenario above, the masses of vector mesons, $\rho$ and $\omega$, should be approximately $2m$, which is well satisfied empirically. Thus, it cannot be overemphasized that the $\pi - \rho$ mass splitting is brought about not by the perturbative color-magnetic interaction between nonrelativistic constituent quarks\[\text{cf.}\] but by the detailed balance between the first and second terms in (2), which is determined exclusively by the demand that the gluonic interaction between current quarks must satisfy chiral $SU(2)_L \times SU(2)_R$ symmetry\[\text{cf.}\]. An important question, which is actively debated nowadays, is which particular nonperturbative gluonic configurations in QCD, e.g. instantons, or abelian monopoles, or

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\[\text{2}\]It is also important to remember that the pion is not a simple nonrelativistic two-body system, but a purely relativistic $\bar{q}q$ system and its Nambu-Goldstone boson nature (zero mass in the chiral limit) cannot be obtained from a nonrelativistic reduction of the second term in (2) used in the Schrödinger equation.

\[\text{3}\]The interaction (2) is color-independent. One can, of course, rewrite this interaction using the Pauli principle in terms of linear combinations of different operators, like $1, color \cdot color, spin \cdot spin, color \cdot color \cdot spin \cdot spin, isospin \cdot isospin,...$ The presence of the
other topological configurations, are intrinsically responsible for the chiral symmetry breaking.

Among other attractive features of the instanton-induced interaction (2) is that it automatically solves the $U(1)_A$ problem, giving a much bigger mass to the pseudoscalar flavor-singlet (in the present 2-flavor formulation that is isosinglet) meson $\eta'$. This is because of the last term in (2). Only this term contributes in a pseudoscalar flavor-singlet quark-antiquark pair. Since this interaction is repulsive, the $\eta'$ becomes heavy, contrary to $\pi$. We note in passing that the color-magnetic interaction cannot explain this big $\eta' - \pi$ mass splitting.

Clearly, the simple Hamiltonian (2) is only some part of a more complicated physical situation. For instance, one definitely needs some additional attractive interaction, e.g. confinement, otherwise the $\eta'$ meson or vector mesons will be unbound, while the octet pseudoscalar mesons are probably not affected by the long-range confining interaction.

Having mentioned all the positive features of the Hamiltonian (2) in the quark-antiquark system, we are now going to discuss its implications in quark-quark systems, i.e. in baryons. What is typically done is a Fierz-rearrangement of the Hamiltonian (3) into diquark $qq$ channels (or, similarly, a Fierz-rearrangement of the generalized NJL Hamiltonian into diquark channels). Then the diquark Hamiltonian is iterated in the $qq$ s-channel, see Fig. 2. The interaction in the scalar $T, J^P = 0,0^+$ diquark turns out to be attractive and it produces a deeply bound scalar diquark.

However, as soon as the Hamiltonian (2) is iterated first in the $qq$ t-channel, see Fig. 3, it implies irreducible (for the $qq$ s-channel) pion- and sigma-exchange interactions between quarks. This statement comes about as a theorem since the iteration of the Hamiltonian (2) in $qq$ t-channel is equivalent to its iteration in $\bar{q}q$ s-channel. Clearly the set of diagrams in Fig. 3 contains all the diagrams of Fig. 2, but in addition it contains many others, and the effect of these additional diagrams is so important that the physics implied by Fig. 2 and Fig. 3 is drastically different.

A simple example of the different physical implications is that Fig. 3 suggests a long-range meson-exchange Yukawa tail, which is crucial for the interaction of quarks, belonging

\textbf{color · color spin · spin} structure in this decomposition does not mean that the interaction (2) becomes similar in its effect to the color-magnetic component of the one gluon exchange interaction, which is explicitly color-dependent. It is like an identity $a = (a - b) + b$ does not mean that effect of $b$ is contained in $a$.

4 In the first calculation [4] the scalar diquark was not bound for $N_c = 3$.

5 We do not show in Fig. 3 a lot of possible chains of bubbles which would correspond to the irreducible two-meson-exchange with crossed meson lines, three-meson-exchange, etc.

6 Sometimes the Hamiltonian (2) is applied in baryons in the framework of the chiral quark-soliton model [13]. In this case quark-quark correlations through the self-consistent chiral mean field and quantization of its rotation take into account some part of the iterations in s-channel of Fig. 2 and do not take into account the t-channel ladders of Fig. 3.
to different nucleons, while if the picture of Fig. 2 were correct the nuclear force would be absent. Another evident difference is that according to the Hamiltonian (2) and Fig. 2 the interaction is absent in flavor-symmetric, $T = 1$, quark pairs $qq$ while the $qq$ interaction of Fig. 3 does not vanish in this case.

Less evident is that even the "short-range" interaction between quarks is crucially modified in Fig. 3 as compared to Fig. 2. We call it the "(anti)screening effect" and illustrate it below.

In order to see it one should avoid the Fierz-rearrangement of (2) into a diquark Hamiltonian. Instead, one can use the initial Hamiltonian (2), but assume that all initial, intermediate and final state $q_iq_j$ wave functions are explicitly antisymmetric.

Consider the first term of (2). In the Nambu-Goldstone mode of chiral symmetry a fermion field has a large dynamical (constituent) mass $m$. Using a $1/m$ expansion, one obtains that to leading order ($m^0$) the first term of (2) leads to a $\delta$-function type attraction in all quark pairs allowed by Pauli principle:

$$-G(\bar{\Psi}\Psi)^2 \implies V(\vec{r}_{ij}) = -2G\delta(\vec{r}_{ij}). \quad (3)$$

The effect of the third term in (2) to the same order is

$$G(\bar{\Psi}\vec{\tau}\Psi)^2 \implies V(\vec{r}_{ij}) = 2G\vec{\tau}_i \cdot \vec{\tau}_j \delta(\vec{r}_{ij}). \quad (4)$$

The potentials (3) and (4), combined together, produce

$$V(\vec{r}_{ij}) = -2G(1 - \vec{\tau}_i \cdot \vec{\tau}_j)\delta(\vec{r}_{ij}). \quad (5)$$

Note that at this order the second and fourth terms of (2) do not contribute. The potential (4) suggests a strong attraction in the isospin-zero quark pair, and no interaction in $T = 1$ quark pairs. Assuming relative angular momentum $L = 0$ within the $T = 0$ quark pair the Pauli principle implies that the spins of the quarks should be antiparallel, $S = 0$. When the theory is sensibly regularized the delta-function attraction is smeared out over the instanton size $r_0$

$$\delta(\vec{r}) \to \frac{1}{4\pi r_0^2} e^{-r/r_0} \frac{1}{r}. \quad (6)$$

This substitution arises from a replacement of the static Green function of the infinitely heavy particle in (2)

$$G_{\mu=\infty}(\vec{x} - \vec{y}) = -r_0^2\delta(\vec{x} - \vec{y}), \quad (7)$$

7I.e. the hyperfine interaction is absent in the $\Delta$-resonance and its excitations. If that were the case, the positive parity state $\Delta(1600)$, which belongs to the $2\hbar\omega$ shell because of its positive parity, would be approximately $\hbar\omega \simeq 500$ MeV above the negative parity pair $\Delta(1620) - \Delta(1700)$.

8The subsequent qualitative discussion will be extended and published in detail elsewhere [16].
by the Green function of a particle with mass $\mu = 1/r_0$

$$G_{\mu=1/r_0}(\vec{x} - \vec{y}) = -\frac{1}{4\pi} \frac{e^{-|\vec{x} - \vec{y}|/r_0}}{|\vec{x} - \vec{y}|}. \quad (8)$$

When the strength of the interaction is big enough, the potential (5)-(6), iterated by solving the semirelativistic Schrödinger equation (i.e. when the kinetic energy operator is taken in a relativistic form) can produce a deeply bound scalar diquark, in agreement with [1,3]. Indeed, when one takes the strength $G = 490 \frac{1}{8N_{c}^2} \text{GeV}^{-2}$, $N_{c} = 3$, with the instanton size $r_0$ between 0.3 and 0.35 fm and the constituent mass $m = 340 - 400 \text{MeV}$ [1,3] one finds a very deeply bound diquark.

In the illustration above we have used a simplified but transparent nonrelativistic picture that adequately reflects in the present case the essential features of a more rigorous Bethe-Salpeter approach.

Sometimes the potential (5) is applied to explain the hyperfine splittings in baryons [17,18]. While it can generate the $\Delta - N$ mass splitting, it fails to explain the lowest levels with positive and negative parity because it does not contain the necessary spin-isospin dependence (1). It will become evident from the discussion below that such an interpretation of the role of instantons in baryons does not survive as soon as the wider class of diagrams in Fig. 3 is considered.

What happens when the first term in (2) is iterated in the $qq$ t-channel? The corresponding amplitude is

$$T_{S}(q^2) = 2G + 2GJ_{S}(q^2)2G + ... = \frac{2G}{1 - 2GJ_{S}(q^2)}, \quad (9)$$

where $J_{S}(q^2)$ is the loop integral (bubble) with the scalar vertex which represents vacuum polarization in the scalar channel. The eq. (9) defines “running amplitude” and a negative sign in the denominator implies its antiscreening behaviour. The expression (9) is known to have a pole at $q^2 = 4m^2$ [12], which can be identified with the exchange by scalar meson $\sigma$ with the mass $\mu_{\sigma} = 2m$ in the chiral limit. The coupling constant of the $\sigma$-meson to constituent quark can be obtained as a residue of (9) at the pole

$$\frac{g_{\sigma q}^2}{q^2 - \mu_{\sigma}^2} = \frac{2G}{1 - 2GJ_{S}(q^2)}. \quad (10)$$

Expanding the $\Psi\sigma\Psi$ vertex in $1/m$, one obtains to leading order ($m^0$) the following well-known sigma-exchange potential

$$V_{\sigma}(r_{ij}) = -\frac{g_{\sigma q}^2}{4\pi} \frac{e^{-\mu_{\sigma}r_{ij}}}{r_{ij}}. \quad (11)$$

The equivalence between the t-channel ladder of bubbles in Fig. 3 beyond the $\sigma$-meson pole in the t-channel and the meson-exchange diagram is achieved only when some form factor $F_{\sigma q}(q^2)$ is inserted into the meson-quark vertex, i.e. the left hand side of eq. (10) should be multiplied with $F_{\sigma q}^2(q^2)$. The form factor is to be normalized $F_{\sigma q}(q^2 = \mu_{\sigma}^2) = 1$. 

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In principle the eq. (10) allows to obtain a functional form for such a form factor. However, it will be very far from reality because the toy model (2) does not contain confinement, which should be important for the interaction between quark and antiquark in the weakly bound system like \( \sigma \)-meson (note that its mass is just at the “continuum threshold” \( 2m \)). In this situation the best way is to rely on our general understanding of the low-energy effective theory. Both constituent quarks and chiral meson fields as well as their couplings make a sense only in the Nambu-Goldstone mode of chiral symmetry. When momentum transfer at the meson - quark vertex exceeds the chiral symmetry breaking scale \( \Lambda_\chi \) (which within all NJL-like models coincides with the regularization scale \( 1/r_0 \)) the effective theory should be cut off. This cut off is accomplished by a form factor in the meson - constituent quark vertex and should be related to the internal structure of both quasiparticles. But in any case the scale parameter in this form factor should be comparable with \( \Lambda_\chi \). At high momenta one can use neither constituent quark nor chiral fields and original quark-gluon degrees of freedom should be used instead.

Approximating this form factor by

\[
F_{\sigma q}^2(q^2) = \frac{\Lambda_\sigma^2 - \mu_\sigma^2}{\Lambda_\sigma^2 - q^2},
\]

instead of the potential (11) we arrive at

\[
V_\sigma(r_{ij}) = -\frac{g_{\sigma q}^2}{4\pi} \left( \frac{e^{-\mu_\sigma r_{ij}}}{r_{ij}} - \frac{e^{-\Lambda_\sigma r_{ij}}}{r_{ij}} \right).
\]

Note, that any functional form of form factor leads to a similar suppression of the potential at short range but it does not influence its long-range part which is determined exclusively by the position of the pole. If one takes a dipole form factor the suppression will be stronger.

What is the fate of the third term in (2), when it is iterated in the \( qq \) t-channel? In this case one obtains the following amplitude

\[
T_{S}^{ab}(q^2) = -\left( 2G - 2GJ_S(q^2)2G + ... \right) \delta_{ab} = -\frac{2G}{1 + 2GJ_S(q^2)} \delta_{ab},
\]

where \( a, b \) are isospin indices. The positive sign in the denominator indicates screening. For instance, at \( q^2 = 4m^2 \) the strength of the interaction is reduced by the factor 2 versus a bare vertex. Still, this suppression of the interaction at low momenta is not realistic, because the toy model (2) does not contain confinement and thus there is only a repulsion in the scalar-isovector quark-antiquark system. When confinement is added in the quark-antiquark pairs, there appear heavy scalar-isovector mesons and the sign of the amplitude (14) becomes opposite at small momenta! This low-momentum amplitude corresponds to the exchange by scalar-isovector mesons between quarks. The corresponding meson-exchange interaction is similar in form to (13), but with an additional factor \( \vec{\tau}_i \cdot \vec{\tau}_j \). The expectation value of the operator \( \vec{\tau}_i \cdot \vec{\tau}_j \) in the scalar diquark is \(-3\) and thus the interaction (3) is stronger by the factor 4 through the interaction (4) in the picture of Fig. 2. In contrast, in the picture of Fig. 3 the contributions from the scalar and scalar-isovector meson exchanges tend to cancel each other.

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Thus we see that the initial interaction is screened. This screening means that the interaction (3) becomes weaken and that its isospin dependence is modified. It is trivial to check that the attraction (13) does not lead to a bound scalar diquark with any reasonable coupling constant, sigma-meson mass, cut-off mass $\Lambda_\sigma$ and constituent quark mass (see discussion in the next chapter). The scalar-isovector meson exchange will further reduce this attraction, though the coupling constant of the scalar-isovector mesons to constituent quarks will be essentially smaller.

Both the scalar-isoscalar exchange and scalar-isovector exchanges between constituent quarks do not contain the flavor-spin dependence (4) which is necessary for baryon spectroscopy.\(^9\)

Now we shall extend our $1/m$ expansion of the Hamiltonian (2) to the next-to-leading order, taking into account terms $m^{-2}$. The first and third terms of (2) will give at this order the spin-orbit forces, as well as some small corrections to the interactions (5) and (6). The second and fourth terms generate, however, a flavor-spin dependent interaction. Consider the second term in (2). It gives both the spin-spin and tensor force components. We ignore below the tensor force as it is irrelevant to our simple discussion and for the $L = 0$ $qq$ pair. Then:

$$- G(\bar{\Psi} i\gamma_5 \vec{\tau} \Psi)^2 \implies V(\vec{r}_{ij}) = 2G \frac{1}{12m^2} \vec{\sigma}_i \cdot \vec{\sigma}_j \vec{\tau}_i \cdot \vec{\tau}_j \Delta \delta(\vec{r}_{ij}). \quad (15)$$

Again, assuming that the instanton has a finite size $r_0$, the potential (13) reads:

$$V(\vec{r}_{ij}) = \frac{2G}{4\pi r_0^2} \frac{1}{12m^2} \vec{\sigma}_i \cdot \vec{\sigma}_j \vec{\tau}_i \cdot \vec{\tau}_j \left( \frac{1}{r_0^2} e^{-r_{ij}/r_0} - 4\pi \delta(\vec{r}_{ij}) \right). \quad (16)$$

Let us now iterate the second term of (4) in the $qq$ t-channel

$$T_P^{ab}(q^2) = \left( 2G + 2G J_P(q^2) 2G + \ldots \right) \delta_{ab} = \frac{2G}{1 - 2G J_P(q^2)} \delta_{ab}, \quad (17)$$

where $J_P(q^2)$ is a bubble with a pseudoscalar vertex (vacuum polarization in the pseudoscalar channel). There is a pole at $q^2 = 0$ in (17) in the chiral limit, which can be identified as a pion-exchange (beyond the chiral limit it is shifted to a physical pion mass $q^2 = \mu^2_\pi$.) The coupling constant of pion to constituent quark can be obtained as a residue of (17) at the pole

$$\frac{g_{\pi qq}^2}{q^2} = \frac{2G}{1 - 2G J_P(q^2)}. \quad (18)$$

\(^9\) We do not state, however, that such an interaction is unimportant. Just the opposite, the $\sigma$-exchange is known to be very important for the medium-range attraction in the NN system and it also contributes to binding nucleon. However it only has a small influence on the splittings via different radial behaviour of baryon wave functions.
Thus near the pole (17) - (18) represents a pion-exchange potential between quarks, which in the chiral limit, $\mu_\pi = 0$, at the order $1/m^2$ (omitting the tensor force component) is:

$$V_\pi(\vec{r}_{ij}) = -\frac{g^2_{\pi q}}{4\pi} \frac{1}{12m^2} \vec{\sigma}_i \cdot \vec{\sigma}_j \vec{\tau}_i \cdot \vec{\tau}_j 4\pi \delta(\vec{r}_{ij}).$$  \hspace{1cm} (19)

The difference between (15-16) and (19) is obvious: the interaction (19) is much stronger.\footnote{One can easily see it from comparison of $\langle \Psi_0 | \Delta \delta(\vec{r}) | \Psi_0 \rangle$ and $\langle \Psi_0 | \delta(\vec{r}) | \Psi_0 \rangle$, where $\Psi_0$ is a zero order function stemming from confinement.} A source of this enhancement is also obvious: near the pole the original bare interaction $2G$ becomes strongly reinforced.\footnote{In essence this antiscreening is some kind of asymptotic freedom: at space-like momenta $q^2 \to -\infty$ the interaction is represented by a bare vertex $2G$, but at $q^2 \to 0$ it becomes infinitely enhanced in the channel with Goldstone boson exchange quantum numbers.} The pion pole is located just near the space-like region and thus strongly influences the quark-quark interaction at not very high momentum transfer.

Again, to retain the equivalence between the t-channel ladder of bubbles in Fig. 3 and the pion-exchange diagram beyond the pole, one must insert a form factor into the $\pi q$ vertex. The effect of this form factor is to smear out the $\delta$-type interaction in (19) over the region $1/\Lambda_{\pi q}$. If this form factor is chosen in the form (12), then one obtains

$$V_\pi(\vec{r}_{ij}) = -\frac{g^2_{\pi q}}{4\pi} \frac{1}{12m^2} \vec{\sigma}_i \cdot \vec{\sigma}_j \vec{\tau}_i \cdot \vec{\tau}_j \frac{e^{-\Lambda_{\pi q} r_{ij}}}{r_{ij}}. \hspace{1cm} (20)$$

The $m^{-2}$ expansion of the fourth term of the Hamiltonian (2) will give a result similar to second term, without, however, isospin-dependent factor and with the opposite sign. Its iteration in the $qq$ t-channel will produce screening effects as it is repulsive in the $\bar{q}q$ s-channel. When, however, this term is combined with an additional attractive interaction, e.g. confinement, it will give $\eta'$. Then the iteration in the $qq$ t-channel will imply $\eta'$-exchange between quarks. The latter interaction is similar to (19), except that the factor $\vec{\tau}_i \cdot \vec{\tau}_j$ is not present and that in this case there appears a Yukawa part of the potential because the mass of $\eta'$ is not zero in the chiral limit.

The discussion above suggests that while for the picture of Fig. 2 the most important interaction is (5) and the interaction (15) is only some very small correction to it, in the case of Fig. 3 the most important interaction in baryons becomes (19) and the one of (13) only plays a modest role for splittings. This is a consequence of an antiscreening. The antiscreening implies that if a typical momentum transfer in the meson-quark vertex (which in $qq$ systems is of the same order as momentum of quarks) is below the chiral symmetry breaking scale, then the original (bare) quark-quark vertex in the pseudoscalar channel is strongly reinforced by the pole which occurs when one iterates it in the t-channel. These pole contributions represent the Goldstone boson exchange interactions between constituent quarks in the Nambu-Goldstone mode of chiral symmetry. Only at a rather high momentum transfer...
transfer (i.e. very far from the poles) there should appear a sensitivity to the original (bare) quark-quark vertex. In the latter case the constituent quarks and chiral fields cannot be used as effective degrees of freedom. So the crucial question is what a typical momentum transfer in the given system is. In the low-lying baryons it is below the chiral symmetry breaking scale thus justifying a use of the effective qq interactions there \[19\].

We hope that the discussion above has been transparent enough to show a dramatic difference between the initial ‘t Hooft interaction, taken literally in \(qq\) system, and its implication after iteration in the t-channel, producing meson exchange between constituent quarks.

In fact, what one needs for the chiral symmetry breaking is a scalar interaction between quarks. Any pairwise gluonic interaction between quarks in the local approximation will necessarily contain the first and second terms of \(2\) with fixed relative strength. This is because of the chiral invariance. Thus all our conclusions are rather general and do not rely necessarily on ‘t Hooft interaction. An important lesson is to see how this nonperturbative gluonic interaction, which induces the dynamical breaking of chiral symmetry, suggests an explanation of both the low-lying mesons and at the same time of baryons and the nuclear force through the effective meson exchange picture in \(qq\) systems. Among the various applications of this idea, will be to work out how the meson-exchange interaction shifts the transition point from the chiral symmetry broken phase to the color-superconductor phase.

We also mention a recent lattice study \[20\] which shows directly that the hyperfine \(\Delta - N\) splitting is mostly due to the meson-exchange interaction between quarks. Another indirect evidence in favor of the picture in Fig. 3 versus that in Fig. 2 is that after cooling (the cooling means that all gluonic configurations, except for instantons, are removed) the \(\Delta - N\) splitting disappears \[21\]. While the cooling does not affect the initial ‘t Hooft interaction between quarks and thus the whole s-channel ladder of Fig. 2 is active, it ruins the t-channel ladder of Fig. 3. The reason is that there are not enough antiquarks in the Fock space after cooling as in quenched approximation they are mostly produced by different gluons, including perturbative ones, attached to valence quark lines (Z graphs).

### III. ARE THE DIQUARK AND NUCLEON BOUND BY THE MESON-EXCHANGE INTERACTION?

We start this section with a short description of the effective meson-exchange interaction model, adjusted to describe baryon spectroscopy within an exact semirelativistic 3-body formulation \[3\]. The Hamiltonian of ref. \[3\] reads:

\[
H = \sum_{i=1}^{3} \sqrt{\vec{p}_i^2 + m_i^2} + \sum_{i<j=1}^{3} V_{ij},
\]

\[
\sum_{i=1}^{3} \vec{p}_i = 0.
\]

Here the relativistic form of the kinetic-energy operator is used, with \(\vec{p}_i\) the 3-momentum and \(m_i\) the masses of the constituent quarks. The dynamical part consists of the quark-quark interaction.
\[ V_{ij} = V_\chi + V_{\text{conf}}. \] (22)

The linear pairwise confining potential
\[ V_{\text{conf}}(r_{ij}) = V_0 + Cr_{ij}, \] (23)
includes both the color-electric string \( Cr_{ij} \) with the color factor absorbed into the string tension \( C \) as well as a constant \( V_0 \), which is large and negative, and thus effectively includes all possible spin- and flavor-independent attractive interactions between quarks, e.g. \( \sigma \)-exchange (13), etc. The flavor- and spin-dependent part of the above Hamiltonian is
\[ V_\chi(r_{ij}) = \left[ \sum_{F=1}^{3} V_\pi(r_{ij})\lambda_i^F\lambda_j^F + \sum_{F=4}^{7} V_K(r_{ij})\lambda_i^F\lambda_j^F \right. \]
\[ + V_\eta(r_{ij})\lambda_i^8\lambda_j^8 + \left. \frac{2}{3} V_{\eta'}(r_{ij}) \right] \vec{\sigma}_i \cdot \vec{\sigma}_j, \] (24)
\[ V_\gamma(r_{ij}) = \frac{g_\gamma^2}{4\pi} \frac{1}{12m_im_j} \left\{ \mu_\gamma^2 e^{-\mu_\gamma r_{ij}} - \Lambda_\gamma^2 e^{-\Lambda_\gamma r_{ij}} \right\}, \] (25)
with \( \mu_\gamma (\gamma = \pi, K, \eta, \eta') \) being the individual phenomenological meson masses, and \( g_\gamma^2/4\pi \) the meson-quark coupling constants.

The constituent mass of the light quarks \( m = m_u = m_d \) was fixed in [6] to a typical value, \( m = 340 \text{ MeV} \), implied by a simple static quark model formula for the nucleon magnetic moment. It is astonishing that the same value has been obtained in a lattice measurement [7]. All other parameters of the above Hamiltonian can be found in ref. [6].

In light quark systems, like \( N \) and \( \Delta \), only the \( \pi \)-like, \( \eta \)-like and \( \eta' \)-like parts of the potential (24) contribute. The \( \pi \)-like exchange interaction is determined by the following matrix elements:
\[-\vec{\tau}_i \cdot \vec{\tau}_j \vec{\sigma}_i \cdot \vec{\sigma}_j = \begin{cases} -1, & \text{if } S_{ij} = 1, \ T_{ij} = 1 \\ -9, & \text{if } S_{ij} = 0, \ T_{ij} = 0 \\ 3, & \text{if } S_{ij} = 1, \ T_{ij} = 0 \\ 3, & \text{if } S_{ij} = 0, \ T_{ij} = 1 \end{cases} \] (26)
while the \( \eta \)- and \( \eta' \)-like exchanges depend only on the spin \( S_{ij} \) of a quark pair.

Combining all \( \pi \), \( \eta \) and \( \eta' \) interactions one finds that the potential (24) is most attractive at short distances in \( S_{ij}, T_{ij} = 0, 0 \) quark pair and essentially less attractive in the \( S_{ij}, T_{ij} = 1, 1 \) diquark system. In other possible color-antitriplet \( q\bar{q} \) pairs it is repulsive.

Applying the Hamiltonian (24) in a color-antitriplet \( qq \) system, one finds a mass \( m_{00} = 744 \text{ MeV} \) for a scalar diquark, \( T, J^P = 0, 0^+ \), and a mass \( m_{11} = 869 \text{ MeV} \) for an axial-vector diquark, \( T, J^P = 1, 1^+ \). In both cases the relative orbital angular momentum is \( L = 0 \), so the total angular momentum coincides with the spin of two quarks. These values are very similar to those obtained recently from the lattice “diquark spectroscopy” [7]. The root-mean-square (r.m.s.) radius of the scalar diquark is 0.354 fm and of the axial-vector one - 0.438 fm. These radii do not include the size of the constituent quark.
It is evident that the confining interaction implies a bound diquark in the sense that there are no asymptotically free constituent quarks. So it is very instructive to compare the mass of the above diquarks with the unphysical two-constituent-mass threshold, \(2m = 680\,\text{MeV}\). The scalar diquark mass is a few tens of MeV above the threshold, which indicates that the meson-exchange part of the interaction, including \(V_0\), does not bind a diquark without the confining interaction \(C_{rij}\). This can also be checked explicitly. To this end we combine the spin- and isospin-dependent interaction (24) with the \(\sigma\)-exchange potential (13) and drop the confining potential (23). The \(\sigma q\) coupling constant is constrained to be equal to the \(\pi q\) one, as suggested by chiral symmetry,

\[
\frac{g_{\sigma q}^2}{4\pi} = \frac{g_{\pi q}^2}{4\pi} = 0.67. \tag{27}
\]

The sigma mass is taken to be \(\mu_\sigma \simeq 2m\), which is implied by the well known result for all NJL-like interactions, \(\mu_\sigma^2 = 4m^2 + \mu_\pi^2\). With these constraints we do not find a bound diquark with any reasonable value for \(\Lambda_\sigma \sim 1\,\text{GeV}\). Only with \(\Lambda_\sigma > 3\,\text{GeV}\) does a weakly bound scalar diquark appear. If one increases the \(\frac{g_{\sigma q}^2}{4\pi}\) coupling constant by a factor 1.5, but keeps the \(\pi q\) coupling constant, then we obtain a bound diquark only at \(\Lambda_\sigma > 1.6\,\text{GeV}\). Thus we conclude that the meson-exchange interaction itself does not bind a diquark.

The next question we address in this section is whether the meson-exchange interaction binds nucleon itself, without confinement. A-priori one cannot exclude the possibility that while the diquark is unbound the three quark system will be bound because of genuine 3-body effects (compare, e.g., the binding energy of tritium and deuteron). Indeed, a full model, including confinement (23), produces a nucleon mass which is below the three constituent mass threshold, \(3m = 1020\,\text{MeV}\). Hence the positive contribution from the rising potential \(C_{rij}\), 631 MeV, is not big compared to the negative contribution from \(3V_0 = -1248\,\text{MeV}\) in combination with the negative contribution of the spin-dependent part of interaction, -750 MeV.

Superficially one could thus conclude that the meson-exchange part of the Hamiltonian could bind nucleon without any support from confinement. However, such an interpretation cannot be taken for two reasons. Firstly, we do not know which part of the negative constant \(V_0\) comes from the \(\sigma\)-exchange, and which - from the genuine color-electric confinement, because the \(Y\)-shape of the gauge-invariant 3-body confining interaction can be approximated by a sum of pairwise potentials only when some additional constant contribution is added. Secondly, a perfect fit of the baryon spectrum with a quality similar to that in ref. [6] can be obtained with a constituent mass smaller than \(m_N/3\). So we have performed a direct calculation of the nucleon, replacing the potential (23) by the \(\sigma\)-exchange potential (13). We have found that for \(\Lambda_\sigma \leq 1800\,\text{MeV}\) the nucleon is unbound and becomes bound at higher values of \(\Lambda_\sigma\). If one increases the \(\frac{g_{\sigma q}^2}{4\pi}\) by a factor 1.5, then the nucleon becomes bound for \(\Lambda_\sigma > 1200\,\text{MeV}\). These results indicate that while the nucleon is unbound with the meson-exchange potential parameters fixed of ref. [6], it could be bound as soon as a spin- and isospin-independent \(\sigma\)-like exchange interaction and/or spin- and isospin-dependent interactions are made stronger, not by a big amount. It is evident that a description of all excited states demands the presence of confinement because all these states are much above the \(3m\) threshold.
IV. IS THERE DIQUARK CLUSTERING IN THE NUCLEON?

We shall use the following set of Jacobi coordinates and a coupling scheme with self-evident notation:

\[ \vec{\rho} = \vec{r}_1 - \vec{r}_2; \quad \vec{\lambda} = \vec{r}_3 - \frac{\vec{r}_1 + \vec{r}_2}{2}, \]

\[ \vec{S}_{12} = \vec{S}_1 + \vec{S}_2; \quad \vec{S} = \vec{S}_{12} + \vec{S}_3, \]

\[ \vec{L} = \vec{L}_\rho + \vec{L}_\lambda, \]

\[ \vec{J} = \vec{L} + \vec{S}, \]

\[ \vec{T}_{12} = \vec{T}_1 + \vec{T}_2; \quad \vec{T} = \vec{T}_{12} + \vec{T}_3. \]

Let \( P_{S_{12}T_{12}} \) be a projector onto a subspace with a given value of spin \( S_{12} \) and isospin \( T_{12} \) of the particles 1 and 2. The probability density for finding particles 1 and 2 in a spin-isospin state \( S_{12}T_{12} \) at a relative distance \( r_{12} \) is given by

\[ g_{S_{12}T_{12}}(r_{12}) = \langle \Psi | P_{S_{12}T_{12}} \delta(\rho - r_{12})|\Psi \rangle, \]

where \( \Psi \) is an antisymmetric 3-body baryon wave function. One can similarly define the probability density for finding the particle 3 at a distance \( r_{12,3} \) from the center of mass of particles 1 and 2

\[ f_{S_{12}T_{12}}(r_{12,3}) = \langle \Psi | P_{S_{12}T_{12}} \delta(\lambda - r_{12,3})|\Psi \rangle. \]

Then one can calculate the corresponding moments

\[ < r_{12}^k > = \int dr_{12} r_{12}^k g_{S_{12}T_{12}}(r_{12}), \]

\[ < r_{12,3}^k > = \int dr_{12,3} r_{12,3}^k f_{S_{12}T_{12}}(r_{12,3}). \]

In Table 1 we present the \( k = 2 \) moments for \( N \) and \( \Delta \) in two cases: (i) full model, (ii) no spin-dependent interaction at all (i.e. only confinement is active).

Comparing the nucleon r.m.s. radius, \( \sqrt{< r_N^2 >} = 0.304 \text{ fm} \), with the radius of a scalar diquark, 0.354 fm, we can deduce the role of genuine 3 body effects - they make the nucleon essentially more compact than the diquark.

The empirical mean square charge radius of the proton, 0.862 \( \text{fm}^2 \), consists of a few contributions: the contribution from the mean square matter radius above, the charge mean square radius of the constituent quark, the meson exchange current contribution \[22,23\], the proton anomalous magnetic moment contribution, etc. The rather small value of the
matter radius, obtained above, is consistent with large contributions from other sources. For instance, the charge radius of the constituent quark should be mainly determined by the $\rho$-meson pole in the time-like region (vector meson dominance) and thus can be expected to be of the order $\sim 0.6$ fm.

The r.m.s. radius of the $\Delta$-resonance, $\sqrt{<r_{\Delta}^2>}$, is $0.390$ fm, is larger than that of the nucleon. This result is easy anticipate since the $\Delta$-resonance wave function does not contain $S_{ij} = T_{ij} = L_{ij} = 0$ components, where the potential (24) is strongly attractive at short range, and thus the size of the $\Delta$-resonance is determined mainly by the weak attraction in the $S_{ij} = T_{ij} = 1$, $L_{ij} = 0$ quark pairs as well as by the confining interaction. The bigger size of $\Delta$ has a well known experimental consequence: the $\Delta \rightarrow N$ electromagnetic form factor falls off faster than the nucleon elastic one.

When the meson-exchange interaction is switched off, the nucleon matter radius becomes larger, $\sqrt{<r_N^2>}$ = $0.442$ fm. This illustrates that there is a soft gap between the scale where chiral physics starts to work and the scale where confinement is important.

The crucial role of three body effects can also be seen from the comparison of the root mean square distance between quarks in the $S_{12}T_{12} = 00$ quark pair in the nucleon, $0.354$ fm, with the same distance in a free scalar diquark, $0.708$ fm. Similarly, the three body effects and the antisymmetrization are responsible for the fact that the root mean square distance in the $S_{12}T_{12} = 00$ quark pair in the nucleon, $0.354$ fm, is similar to that one in the $S_{12}T_{12} = 11$ subsystem in the nucleon, $0.387$ fm, while the potential is very different in both cases. A comparison of the two numbers above gives an idea about how unimportant clustering is in the nucleon. It can also be seen from Fig. 4 and Fig. 5 where we show probability density distributions.

With a pure static “right triangle” $3q$ configuration the relation between $<r_{12}^2>$ and $<r_{12,3}^2>$ would be $<r_{12,3}^2> = \frac{3}{4} <r_{12}^2>$. This relation is almost exactly satisfied with the $\Delta$ wave function or with the $3q$ wave function when the meson-exchange interaction is switched off. In the nucleon wave function there is a deviation from this relation but it is not large. We thus conclude that there is no appreciable clustering in the nucleon.

What is the physical reason for the absence of a significant clustering? The answer is that there are genuine 3-body effects and the fermi-nature of quarks do not support clustering. Indeed, if the quarks, say, with numbers 1 and 2 form a pair $S_{12}T_{12} = 00$, the antisymmetry of the wave function suggests that there are at the same time pairs with quantum numbers $S_{13}T_{13} = 00$ or $S_{23}T_{23} = 00$ (along with other quantum numbers). Thus a strong attraction acts simultaneously in all quark pairs which makes the nucleon compact but not clustered much. Only a much stronger and “sharper” interaction in the $ST = 00$ diquark would lead to an appreciable clustering, but at a cost that $\Delta - N$ splitting will become enormous.

**V. SUMMARY**

Here we summarize our main conclusions.

1. The nonperturbative gluonic interaction between quarks, e.g. instanton-induced one, which is responsible for the dynamical breaking of chiral symmetry in QCD and thus explains the $\pi - \rho$ mass splitting, iterated in the $qq$ t-channel implies a meson-exchange picture between constituent quarks, and through the latter also explains baryons and nuclear force.
This is a simple consequence of crossing symmetry: if one obtains pion as a solution of the Bethe-Salpeter equation in the quark-antiquark s-channel, then one inevitably obtains a pion-exchange in the quark-quark systems as a result of iterations in the $qq$ t-channel.

2. Due to (anti)screening effects the implications of this nonperturbative gluonic interaction in $qq$ systems are drastically different when it is iterated only in the s-channel as compared to a more general case, when it is first iterated in the t-channel, leading to a meson exchange, and only after that iterated in the s-channel.

3. The effective meson-exchange interaction in $qq$ systems does not bind diquarks without an additional confining force and does not induce any appreciable clustering in the nucleon.

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TABLES

TABLE I. Relative distances in $N$ and $\Delta$

| system       | $(S_{12}, T_{12})$ | probability | $\langle r_{12}^2 \rangle$, fm$^2$ | $\langle r_{12,3}^2 \rangle$, fm$^2$ |
|--------------|--------------------|-------------|-----------------------------------|-----------------------------------|
| Nucleon      | (0,0)              | 0.498       | 0.125                             | 0.114                             |
| Nucleon      | (1,1)              | 0.498       | 0.150                             | 0.094                             |
| Delta        | (1,1)              | 1.000       | 0.454                             | 0.342                             |
| Conf. only   | (0,0) or (1,1)     |             | 0.585                             | 0.441                             |
FIG. 1. The s-channel ladder in $q\bar{q}$ system.

FIG. 2. The s-channel ladder in $qq$ system.
FIG. 3. The s- and t-channel ladders in $qq$ system.

FIG. 4. The probability distributions (for a definition see the text). Solid line - $g_{00}(r_{12})$ for $N$; dotted line - $g_{11}(r_{12})$ for $N$; dashed line - $g_{11}(r_{12})$ for $\Delta$; long dashed line - $g_{00}(r_{12}) = g_{11}(r_{12})$ for $N$ when the spin-dependent interaction is switched off.
FIG. 5. The probability distributions (for a definition see the text). Solid line - $f_{00}(r_{12,3})$ for $N$; dotted line - $f_{11}(r_{12,3})$ for $N$; dashed line - $f_{11}(r_{12,3})$ for $\Delta$; long dashed line - $f_{00}(r_{12,3}) = f_{11}(r_{12,3})$ for $N$ when the spin-dependent interaction is switched off.