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ЯДЕРНЫЕ СИЛЫ И КВАНТОВАЯ ХРОМОДИНАМИКА

Обсуждается механизм возникновения ядерных сил. Предполагается, что на малых расстояниях между нуклонами происходит обмен цветными кварками в согласии с основными положениями квантовой хромодинамики. Показано, что помимо привычных двухчастичных сил, в системе трех и более нуклонов должны существовать истинно трехчастичные силы.

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1. Introduction

For many years of the nuclear physics development it was considered that nucleons and mesons were point-like structureless elementary particles. In works of Heisenberg, Wigner, Majorana, Yukawa and others many of the nucleus properties were explained proceeding from two-body forces only acting between point-like particles. However at present it is impossible to ignore the fact that nucleons are non-elementary and according to the quark model are composed of three quarks, and mesons - of a quark-antiquark pair. The experimental data on the asymptotic behaviour of the deuteron electromagnetic form-factor[1] interpreted with the help of quark counting rules[2,3] were the first direct indication of the possible display of quark structure in the atomic nuclei.

On the other hand it turned out that by limiting oneself to two-nucleon potentials only it was impossible to simultaneously describe correctly the binding energy and charge form-factor of light nuclei, particularly, the mean-square radius and the minimum position in the form-factors of the nuclei with A = 3 (^{3}H, ^{3}He) and A = 4 (^{4}He). To improve this state a "true" three-body potential should be added to an ordinary two-body Hamiltonian. Some aspects of introducing many-body forces on the basis of the quark model are discussed by Robson [4].

In the present work it is considered that nuclear forces in light nuclei arise mainly at the expense of coloured quark exchange at small distances between nucleons. As it is known, in the quantum chromodynamics (QCD) the colour group is a second rank group SU(3)c. Therefore it makes possible the existence of two types of exchange forces - ordinary two-body P_{ij} and three-body P_{ijk}. The obtained in the work ratio between the binding energies of light nuclei (see below formula (11)) does not contradict the experimental data and testifies that the introduced in this way two- and three-body exchange forces are realistic.

2. Exchange Forces in QCD Perturbation Theory

At present it is accepted to consider that the nucleon-nucleon forces may be understood on the basis of quark-gluon processes proceeding from the QCD fundamental principles - non-Abelian gauge theory of strong interactions. The strict application of quantum chromodynamics to the problem of nuclear forces is not possible for the time being, since the problem of confinement, i.e. absence of free colour states is not solved yet. However the QCD simplified model allows to obtain a number of results which apparently may be more well-grounded due to the further development of the theory.

According to QCD, the observed hadrons, in particular nucleons are the group SU(3)c scalars or, as it is customary to say, nucleons are colourless. Since a gluon belongs to the group SU(3)c octet representation, the simplest diagram of one gluon exchange between two nucleons turns into zero (Fig. 1a).

This may also be proved by a direct calculation of colour factors in accord with QCD rules. Indeed (with the accuracy up to a multiplicative constant) for the Fig. 1a diagram we have:

$$\frac{1}{6^2}\varepsilon_{\alpha\beta\gamma}\delta_{\alpha\alpha'}\delta_{\beta\beta'}(\lambda_a)_{\gamma\gamma'}\varepsilon_{\alpha'\beta'\gamma'}\varepsilon_{\delta\mu\nu}(\lambda_a)_{\delta\delta'}\delta_{\mu\mu'}\delta_{\nu\nu'}\varepsilon_{\delta'\mu'\nu'} = \frac{1}{9}\delta_{\gamma\gamma'}(\lambda_a)_{\gamma\gamma'}\,\delta_{\delta\delta'}(\lambda_a)_{\delta\delta'} = \frac{1}{9}(Sp\,\lambda_a)^2 = 0 \tag{1}$$

Here $\alpha, \beta, \ldots, \mu', \nu'$ are the quark colour indices adopting values 1,2,3; $\frac{1}{\sqrt{6}}\varepsilon_{\alpha\beta\gamma}$ is the colour wave function of a colourless nucleon; λ_a $(a=1,\ldots,8)$ are the Gell-Mann 3×3 matrices. Let's show that the Fig. 1b diagram with quark exchange between nucleons is already not equal to zero. Indeed for the diagram 1b we have:

$$\frac{1}{6^2}\varepsilon_{\alpha\beta\gamma}\delta_{\alpha\alpha'}\delta_{\beta\beta'}(\lambda_a)_{\gamma\gamma'}\varepsilon_{\alpha'\beta'\gamma'}\varepsilon_{\delta\mu\nu}(\lambda_a)_{\delta\delta'}\delta_{\mu\mu'}\delta_{\nu\nu'}\varepsilon_{\delta'\mu'\nu'} = \frac{1}{9}\delta_{\gamma\gamma'}(\lambda_a)_{\gamma\delta'}\,\delta_{\delta\delta'}(\lambda_a)_{\delta\gamma'} = \frac{1}{9}(\lambda_a)_{\gamma\delta}(\lambda_a)_{\delta\gamma} = \frac{4}{9} \neq 0 \tag{2}$$

Here at passing to the last line the following identity was used:

$$(\lambda_a)_{\alpha\beta}(\lambda_a)_{\gamma\delta} = \frac{1}{2}\left(\delta_{\alpha\delta}\delta_{\gamma\beta} - \frac{1}{3}\delta_{\alpha\beta}\delta_{\gamma\delta}\right) \qquad (3)$$

Without going into the analysis of other possible diagrams of perturbation theory in the two-nucleon system, let's pass to diagrams describing interactions of three nucleons. With the help of the Fig. 1b diagram, by adding to it one more nucleon which exchanges analogous quarks with two other nucleons (such a diagram is convenient to draw on the cylinder surface) it is possible to describe a three-nucleon interaction by means of ordinary two-body potentials of the form $V_{12} + V_{13} + V_{23}$. A "true" three-body force between three nucleons, described by the V_{123}-form potential may arise due to three-gluon vertex. Let's consider the system of three nucleons connected by the three-gluon vertex. The Fig. 2a diagram, in which there is no quark exchange between nucleons, is equal to zero. One can make sure of it by both direct calculation and reasoning, analogous to the proof of the Fig. 1a diagram being equal to zero Substituting colour factors in the diagram 2a we have

$$\frac{1}{6^3}\varepsilon_{\alpha\beta\gamma}\delta_{\alpha\alpha'}\delta_{\beta\beta'}\varepsilon_{\alpha'\beta'\gamma'}\varepsilon_{\delta\mu\nu}\delta_{\mu\mu'}\delta_{\nu\nu'}\varepsilon_{\delta'\mu'\nu'}\varepsilon_{\sigma\tau\rho}\delta_{\tau\tau'}\delta_{\rho\rho'}\varepsilon_{\sigma'\tau'\rho'}$$
$$\times\ f_{abc}(\lambda_a)_{\gamma\gamma'}(\lambda_b)_{\delta\delta'}(\lambda_c)_{\sigma\sigma'} = \qquad (4)$$
$$= \frac{1}{6^2} f_{abc}\ \mathrm{Sp}(\lambda_a)\ \mathrm{Sp}(\lambda_b)\ \mathrm{Sp}(\lambda_c) = 0$$

Here f_{abc} is the structure constant of the group SU(3)c.

But the diagram 2b with a subsequent quark exchange between three nucleons is different from zero. Indeed substituting the colour factors in the diagram 2b we have:

$$\frac{1}{6^3}\varepsilon_{\alpha\beta\gamma}\delta_{\alpha\alpha'}\delta_{\beta\beta'}\varepsilon_{\alpha'\beta'\gamma'}\varepsilon_{\delta\mu\nu}\delta_{\mu\mu'}\delta_{\nu\nu'}\varepsilon_{\delta'\mu'\nu'}\varepsilon_{\sigma\tau\rho}\delta_{\tau\tau'}\delta_{\rho\rho'}\varepsilon_{\sigma'\tau'\rho'}\ \times$$
$$\times\ f_{abc}(\lambda_a)_{\delta\gamma'}(\lambda_b)_{\gamma\sigma'}(\lambda_c)_{\sigma\delta'} = \qquad (5)$$
$$= -\frac{16i}{9} \neq 0$$

Here we have made use of the identity

$$f_{abc}(\lambda_a)_{\alpha\alpha'}(\lambda_b)_{\beta\beta'}(\lambda_c)_{\gamma\gamma'} = 2i\left(\delta_{\alpha\gamma'}\delta_{\beta\alpha'}\delta_{\gamma\beta'} - \delta_{\alpha\beta'}\delta_{\beta\gamma'}\delta_{\gamma\alpha'}\right)$$

Thus the Fig. 2b diagram leads to a true three-body V_{123}-type potential.

We didn't discuss the contribution of many other diagrams. But the examples discussed showed that, according to QCD, interactions of <u>exchange</u> character are possible between colourless hadrons.

Certainly in the perturbation theory higher orders there are non-zero diagrams which do not contain quark exchange. Such diagrams, leading to the van der Waals forces at the expense of colour polarization in the intermediate states, have been evaluated in the work [6] and apparently they are smaller as compared with exchange forces. In the next section the exchange forces are introduced for the general group reasons without application of perturbation theory.

3. Two-Quark and Three-Quark Exchange Forces

In the present section the discussion of exchange forces in QCD is carried on with the help of a method which is the generalization of the Dirac method who was the first to introduce spin exchange forces in atomic physics [7].

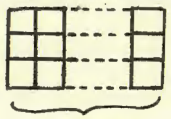
According to quantum chromodynamics quarks (see e.g. review [5]) have a colour which is the source (charge) of a strong interaction between quarks carried by gluon fields. The quark and gluon fields exactly satisfy the inner SU(3)c symmetry called colour symmetry. Since group SU(3)c is a second rank one, its irreducible representations are characterized by two Casimir operators - a quadratic and a cubic. Therefore the group allows exchange interactions of two types only - two- and three-body. From the SU(3)c - invariance it follows that the quark system energy operator, which is in the singlet colour state, should have the form

$$E = a \sum_{i<j}^{N} P_{ij} + b \sum_{i<j<k}^{N} P_{ijk} \qquad (6)$$

where P_{ij} and P_{ijk} are the permutation operators of two and three quarks respectively.

Remind that nucleon is a colourless three-quark combination, the colour wave function of which is described by the Young diagram of the form $[1^3]$. The colourless combination of six quarks has the Young diagram $[2^3]$. Likewise diagrams $[3^3]$ and $[4^3]$ correspond to 9- and 12-quark colourless states. Gene-

rally an arbitrary colourless nucleus with the number of nucleons equal to A, consisting of 3A quarks is described by the Young diagram having 3 lines and A columns



A times

The matrix elements of two- and three-body permutation operators for the states with the Young diagram $[A^3]$ are equal to (see Appendix):

$$\langle [A^3] | \sum_{i<j} P_{ij} | [A^3] \rangle = \frac{3}{2} A (A-3) \qquad (7)$$

$$\langle [A^3] | \sum_{i<j<k} P_{ijk} | [A^3] \rangle = \frac{1}{2} A (A^2-9A+10) \qquad (8)$$

Using (6), (7) and (8) it is easy to extract energies of multi-quark colourless states through parameter a and b:

$$\begin{aligned} E_{3q} &= -3a + b \\ E_{6q} &= -3a - 4b \\ E_{9q} &= -12b \\ E_{12q} &= 6a - 20b \end{aligned} \qquad (9)$$

The binding energy of light nuclei - deuteron, ^3He and ^4He may be obtained from here by means of subtracting the binding energy of three-qiark combinations corresponding to colourless nucleons:

$$E(^2H) = E_{6q} - 2E_{3q} = 3a - 6b$$
$$E(^3He) = E_{9q} - 3E_{3q} = 9a - 15b \qquad (10)$$
$$E(^4He) = E_{12q} - 4E_{3q} = 18a - 24b$$

Excluding the two constants a and b from these three equations we shall obtain the relation between binding energies of deuteron, ^{3}He and ^{4}He:

$$E(^4He) = 4E(^3He) - 6E(^2H) \qquad (11)$$

The relation (11) is valid in the limits of exact spin-isospin SU(4)-symmetry of Wigner when it is ignored by spin and isospin (i.e. electromagnetic) interactions. Therefore when comparing the relation (11) with experimental data the mean value of the binding energy of ^{3}He and tritium ^{3}H should be substituted for E(^{3}He):

$$E(^3He) \to \frac{1}{2}\left[E(^3He) + E(^3H)\right] = \frac{1}{2}(7{,}72 + 8{,}49)\,\text{Мэв} = 8{,}2\,\text{Мэв}$$

and for E(^{2}H) - the mean value of the binding energy of deuteron and virtual S_0 - the state of proton and neutron (equal to $\approx$ 0):

$$E(^2H) \to \frac{1}{2}\left[E(^2H) + E(^1S_0)\right] = \frac{1}{2}(2{,}21 + 0)\,\text{Мэв} = 1{,}1\,\text{Мэв}$$

Helium-4 is the singlet of the Wigner group SU(4), therefore the exact value of E(^{4}He) = 28.30 MeV may be substituted in the relation (6). With the account of the abovementioned the relation (11) satisfies the experimental data fairly well (28.3 MeV $\approx$ 25.8 MeV).

It is interesting to note that Robson's intuitive considerations [4] on the character of three-nucleon forces lead to the same relation (6). Indeed, following the work [4], let's mark the distribution of two-nucleon forces through b', and that of three-nucleon forces - through Δ. Then, ignoring the Coulomb forces, with the help of a simple combinatorial calculation we shall obtain the expressions for the binding energies of light nuclei:

$$E(^2H) = b'$$
$$E(^3He) = E(^3H) = \Delta + 3b' \qquad (12)$$
$$E(^4He) = 4\Delta + 6b'$$

It is easy to notice that the exclusion of parameters b' and Δ leads again to the relation (11). The comparison of (7) with (5) shows that the thus introduced nucleon parameters are simply connected with quark parameters a and b: Δ = 3b, b' = 3a-6b.

Up to now we have dealt with ordinary nuclei consisting of non-strange quarks u and d. The maximum number of non-strange quarks which can simultaneously be in the state with a zero orbital momentum (S state) is equal to 2(flavor) $\times$ $\times$ 2(spin) $\times$ 3(colour) = 12 which corresponds to nucleus ^{4}He composed of two protons and two neutrons. From a quark structure viewpoint the helium nucleus has a configuration 6u6d. The strange quark s introduced the number of quarks which can simultaneously be in the S state will make 18 [3(flavor) $\times$

$2(\text{spin}) \times 3(\text{colour}) = 18]$. It is clear that this will be a state composed of 6u6d6s quarks. From a flavor viewpoint this combination has the same quark composition as a hypernucleus composed of six Λ-hyperons since 6u6d6s = 6(uds) = 6Λ. However this will not be the lowest energy level of the system since, owing to the Pauli principle, only two Λ-hyperons can simultaneously be in the state S. Therefore at such a grouping of quarks it is impossible to construct a nucleus in which all the quarks are in the S state.

However, considering the quark structure of the proton P = uud, neutron n = udd and Ω-hyperon Ω = sss, the combination of 18 quarks in the S state can be expanded into baryons in the following way 6u6d6s = 2p2n2Ω. One can imagine such a hypernucleus, with strangeness 6 and zero electric charge, as a nucleus of ^{4}He having joined two Ω-hyperons. It is theoretically possible that there exists a simpler hypernucleus composed of only one Ω-hyperon, two protons and two neutrons.

Since all the baryons in these hypernuclei are in the S state, one can expect a large contribution of three-body forces which leads to a relatively large binding energy for these hypernuclei. The search of such nuclei, for example, in cosmic rays will offer an important and interesting problem.

Conclusion

The discussed in the present work two-quark and three-quark exchange forces apparently make the main contribution

at small distances between nucleons. As for large distances, it occurs that the π-meson exchange mechanism remains valid. From the quark model viewpoint the exchange of π-mesons (and other mesons) between colourless hadrons is possible in higher approximations of perturbation theory at the expense of presence of virtual quark-antiquark loops in interacting nucleons. And since the nuclei ^{3}H, ^{3}He and ^{4}He are closely bound systems with a rather small mean-square radius, it may turn out that their binding energy, on the whole, is due to the discussed here quark exchange mechanism.

It is very difficult to experimentally separate the three-nucleon force effect from the two-body force extrapolated out of energy surface. However the data on the light nuclei binding energy and mean-square radii obtained from the form factor measurements more and more testify to the real existence of three-body forces. The present work shows that the quark model of nuclei based on the simplified interpretation of quantum chromodynamics results naturally in the necessity of taking into account the three-quark exchange forces leading to effective three-nucleon forces.

The new approach to the problem of nuclear forces on the basis of QCD simplified model allows to hope that the secret of the main paradigm of modern nuclear physics - the Yukawa π-meson exchange, will be understood in the nearest future just on this way.

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Appendix

In the present appendix the properties are studied of permutation operators of two and three particles P_{ij} and P_{ijk} in the general case of the symmetry U(n) group.

As it is known, group U(n) has $n \times n$ generators A^α_β ($\alpha, \beta = 1, 2, \ldots n$). In the lowest, i.e. fundamental or quark representation the generators satisfy the relation

$$A^\alpha_\beta(i)\, A^\gamma_\delta(i) = \delta^\gamma_\beta\, A^\alpha_\delta(i) \tag{A.1}$$

For the system of N particles generators have the form $A^\alpha_\beta = \sum_{i=1}^N A^\alpha_\beta(i)$ and satisfy the permutation relations of group U(n):

$$[A^\alpha_\beta, A^\gamma_\delta] = \delta^\gamma_\beta A^\alpha_\delta - \delta^\alpha_\delta A^\gamma_\beta \tag{A.2}$$

The arbitrary irreducible representation of group U(n) is characterized by the Young diagram ($f_1, f_2, \ldots, f_n$) where f_i is the number of cells in the i-th line of the Young diagram. On the other hand, one may introduce n Casimir operators the eigenvalues of which are expressed through the partition numbers of the Young diagram:

$$\langle C_1^{U(n)} \rangle = \langle A^\alpha_\alpha \rangle = f_1 + f_2 + \ldots + f_n = N$$

$$\langle C_2^{U(n)} \rangle = \langle A^\alpha_\beta A^\beta_\alpha \rangle = \sum_{k=1}^n f_k^2 + \sum_{k=1}^n (n+1-2k) f_k \tag{A.3}$$

$$\vdots$$

$$\langle C_p^{U(n)} \rangle = \langle A^\alpha_\beta A^\beta_\gamma \ldots A^\mu_\alpha \rangle = \sum_{k=1}^n (l_k)^p \prod_{m=1}^{n}{}' \frac{1 + l_m - l_n}{l_m - l_n}$$

In the last formula $l_k = f_k + n - k$, the prime by the product symbol means that $m \neq k$, and the number P takes values 1,2,...n [10].

By the definition (see e.g. [8])the two- and three-body permutation operators are expressed through the group generators in the following way:

$$P_{ij} = A^\alpha_\beta(i)\, A^\beta_\alpha(j) \tag{A.4}$$

$$P_{ijk} = A^\alpha_\beta(i)\, A^\beta_\gamma(j)\, A^\gamma_\alpha(k) \tag{A.5}$$

Corresponding sums of two- and three-body permutation operators are expressed through the Casimir operators:

$$\sum_{i \neq j}^N P_{ij} = C_2^{U(n)} - nN \tag{A.6}$$

$$\sum_{i \neq j \neq k}^N P_{ijk} = C_3^{U(n)} - 2n\, C_2^{U(n)} - (n^2+1)N - N^2 \tag{A.7}$$

The first of these relations is known and easy to prove. Let's give a short derivation of the second one. By the definition of the third rank Casimir operator we have:

$$C_3^{U(n)} = A^\alpha_\beta A^\beta_\gamma A^\gamma_\alpha =$$

$$= \sum_{i=1}^N A^\alpha_\beta(i) \sum_{j=1}^N A^\beta_\gamma(j) \sum_{k=1}^N A^\gamma_\alpha(k) =$$

$$= \sum_{i,j,k}^N A^\alpha_\beta(i) A^\beta_\gamma(j) A^\gamma_\alpha(k) =$$

$$= \sum_{i \neq j \neq k}^N A^\alpha_\beta(i) A^\beta_\gamma(j) A^\gamma_\alpha(k) + \sum_{i=1}^N A^\alpha_\beta(i) A^\beta_\gamma(i) A^\gamma_\alpha(i) + \tag{A.8}$$

$$+\sum_{i\neq j}^{N} A^{\alpha}_{\beta}(i)A^{\beta}_{\gamma}(i)A^{\gamma}_{\alpha}(j) + \sum_{i\neq j}^{N} A^{\alpha}_{\beta}(i)A^{\beta}_{\gamma}(j)A^{\gamma}_{\alpha}(i) + \sum_{i\neq j}^{N} A^{\alpha}_{\beta}(j)A^{\beta}_{\gamma}(i)A^{\gamma}_{\alpha}(i) =$$

$$= \sum_{i\neq j\neq k} P_{ijk} + n^2 N + n\sum_{i\neq j} P_{ij} + N(N-1) + n\sum_{i\neq j} P_{ij}$$

Hence

$$\sum_{P}{}' P_{ijk} = C_3^{U(n)} - 2n\sum_{i\neq j} P_{ij} - (n^2-1)N - N^2 \qquad (A.9)$$

Substituting here the value $\sum_{i\neq j} P_{ij}$ through the Casimir second rank operator we shall obtain the required equation (A.7). The operator eigenvalue $\sum_{i\neq j\neq k} P_{ijk}$ in the case of the group U(3) is easy to obtain from here considering that

$$\langle C_2^{U(3)} \rangle = f_1^2 + f_2^2 + f_3^2 + 2f_1 - 2f_3 \qquad (A.10)$$

$$\langle C_3^{U(3)} \rangle = f_1^3 + f_2^3 + f_3^3 + 4f_1^2 + f_2^2 - 2f_3^2 - - (f_1 f_2 + f_1 f_3 + f_2 f_3) - 4f_1 - 2f_2 - 2f_3 \qquad (A.11)$$

Therefore in the case of group U(3) we shall obtain (see also [9]):

$$\sum_{i\neq j} P_{ij} = f_1(f_1-1) + f_2(f_2-3) + f_3(f_3-5) \qquad (A.12)$$

$$\sum_{i\neq j\neq k} P_{ijk} = f_1(f_1-1)(f_1-2) + f_2(f_2-2)(f_2-4) + + f_3(f_3-4)(f_3-5) - 3(f_1 f_2 + f_1 f_3 + f_2 f_3) \qquad (A.13)$$

In the main text of the paper $f_1 = f_2 = f_3 = A$. Taking into account also that $\sum_{i\neq j} = 2\sum_{i<j}$ and $\sum_{i\neq j\neq k} = 6\sum_{i<j<k}$

we shall obtain the formulae(7) and (8). To make it complete let's note that the total exchange energy of the quark system 3A with account of two- and three-body forces is still easier expressed directly through the Casimir operators

$$E = \alpha\, C_2^{U(3)} + \beta\, C_3^{U(3)} \qquad (A.14)$$

or

$$\langle E \rangle = 3\alpha A^2 + 3\beta A^3 \qquad (A.15)$$

which results in the same basic relation (11).

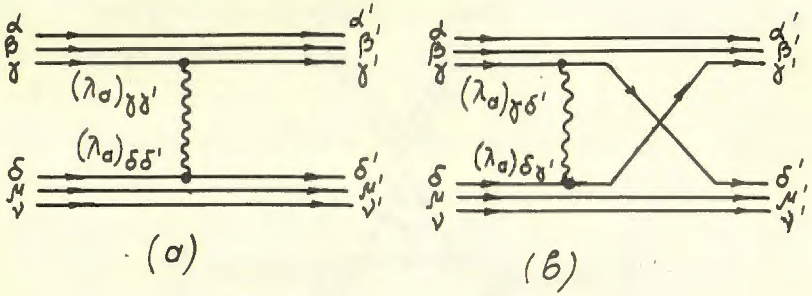

Fig. 1 Diagram of one-gluon exchange between two colourless hadrons:

a) without quark exchange (= 0)

b) with quark exchange (≠ 0)

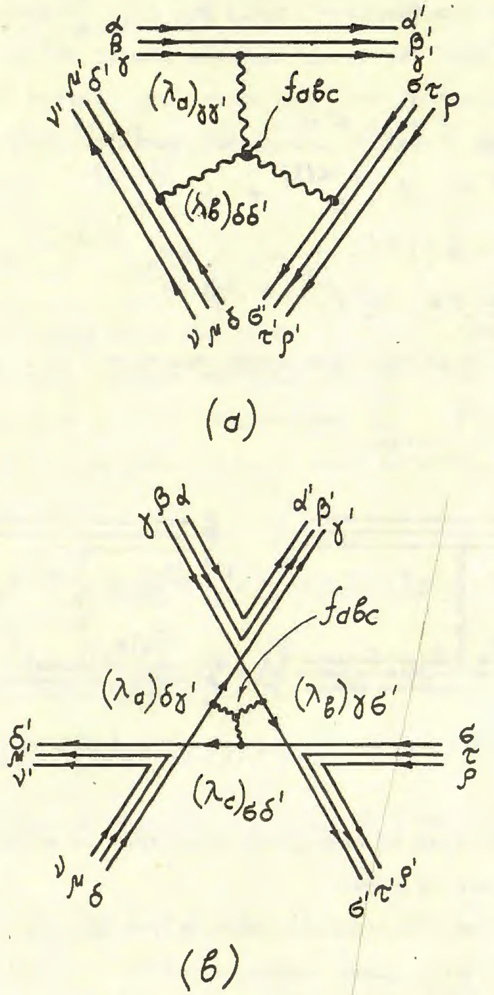

Fig. 2 Three-gluon vertex in the three-nucleon system:
a) without quark exchange (= 0)
b) diagram illustrating the three-quark exchange at the presence of three-gluon vertex ($\neq 0$).

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