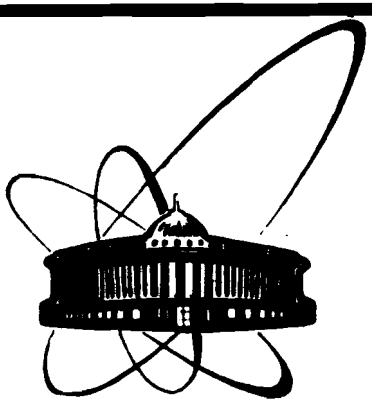

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PAULI-QUENCHING FOR HADRONS IN NUCLEAR MATTER – A QUARK SUBSTRUCTURE EFFECT

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1. Introduction

At the present there is a growing interest to understand the properties of nuclear matter on the basis of the underlying quark substructure. Whereas now a first principle QCD-approach to this problem cannot be realised, a semiphenomenological quark potential approach can be successfully applied to work out the description of hadronic properties within a quark picture. Nonrelativistic quark models have been proven remarkably useful in describing the hadron spectroscopy [1]. Many efforts have been made to derive the hadron-hadron interaction from say the six-quark problem. Phase shifts obtained from a non-relativistic quark potential model give a good fit to the nucleon-nucleon and meson-meson scattering data [2,3,7,8,12].

Another interesting problem is the investigation of nuclear matter as a many quark system at finite temperature and density. As a consequence of their quark substructure the nucleons are affected by the surrounding nuclear medium. In contrast to the few-quark problem, where we have to solve the Schrödinger equation, a quantum statistical approach is needed to treat the many-quark system at finite temperature. Because of the confinement property of the quark interaction potential, this quantum statistical approach must be modified if compared with usual classical many-particle systems.

To formulate the Hamiltonian we consider non-relativistic massive quarks so that the kinetic energy is given by

$$KE = \sum_{i=1}^{N} \left(m + \frac{p_i^2}{2m} \right) . \qquad (1)$$

The potential energy $PE(r_1 \ldots r_N)$ is constructed in the following

way [5]: The configuration $(\underline{r}_1 \ldots \underline{r}_N)$ is decomposed into color-neutral clusters $q\bar{q}$ or qqq, respectively. The potential

$$V_{ij} = \frac{m\omega^2}{2} (r_i - r_j)^2 \quad , \tag{2}$$

is assumed to act only within these color-neutral clusters (saturation property of the interaction). Within all possible decompositions of the quark configurations one has to take the cluster configuration with the minimum potential energy, this minimum value of the potential energy will be denoted by PE.
The Hamiltonian is then given by

$$H = KE + PE \quad . \tag{3}$$

Of course, this Hamiltonian is able to describe isolated hadrons where the quark interaction is confined within the color-neutral hadronic cluster. With respect to the two-nucleon problem [3], color van-der-Waals forces do not arise because of the saturation property of the quark interaction. A massive quark matter phase can be described where the potential energy is given by the distribution function of the next neighbours [4].

We consider nuclear matter as the hadronised phase where the interaction strings are confined within the nucleons and string flips like in the quark matter phase are not likely to occur. However, the color-neutral three-quark cluster is influenced by the surrounding clusters by reason of the Pauli principle what demands the antisymmetrisation of the hadronic quark wave functions. The corresponding shift of the nucleon energy which may be considered as the self-energy of the three-quark cluster should contribute to the binding energy of nuclear matter. It is the aim of this work to evaluate this self-energy contribution due to Pauli blocking and to compare it with empirical values for the binding energy of nuclear matter.

Within a Green function approach [5] the lowest order diagram with respect to the density gives the Pauli shift of the three-quark cluster

$$\Delta E_n^{Pauli} = \sum_{n'} \Delta E_{nn'}^{Pauli} \; f_3(E_{n'}) \quad ;$$

$$\Delta E_{nn'}^{Pauli} = 3 \sum_{1\ldots6} \Psi_n^*(123)\Psi_{n'}^*(456)\{KE - E_n - E_{n'}\}\{\Psi_n(126)\Psi_{n'}(453) - \Psi_n(453)\Psi_{n'}(126)\}. \tag{4}$$

This Pauli blocking shift has already been evaluated for finite temperatures and densities of the nuclear environment and leads to temperature and density dependent nucleonic properties, such as the effective nucleon mass [6], effective nucleonic radii [13] and corresponds to the hard-core part of the effective Skyrme interaction for nuclear matter [5].

However, at zero temperature the Pauli quenching shift (4) obtained within a quantum statistical treatment of the completely hadronised quark plasma may be interpreted as a contribution due to an appropriately chosen antisymmetrisation of the six-quark wave function $\phi_{nn'}(1\ldots6)$ of the two-nucleon problem. In the present work we want to show this correspondence in detail thus contacting few-body approaches which deal with the problem of effective NN-interactions on the quark level using the Resonating group method [7-9].

2. Antisymmetrisation of the Two-Nucleon Wave Function $\phi_{nn'}(1\ldots6)$

We want to represent the six-quark wave function $\phi_{nn'}(1\ldots6)$ as a product of two nucleonic wavefunctions that behaves antisymmetrically with respect to each exchange of quantum numbers belonging to quarks $(P_{ij};\ i=1,2,3\ \ j=4,5,6)$ or to nucleons $(P_{nn'})$ thus fulfilling the Pauli principle on the nucleonic as well as on the quark level. Following this prescription and considering only the two-nucleon channel, all those permutations leading to color non-singlet clusters have to be excluded and we obtain

$$\phi_{nn'}(1\ldots6) = (1 - \sum_{i=1}^{3} P_{i,i+3})(1 - P_{nn'})\,\Psi_n(123)\Psi_{n'}(456) \quad , \tag{5}$$

where the numbers i=1...6 stand for the momentum, spin, flavor and color indices of the ith quark and n denotes the center-of-mass momentum P as well as one of the spin-isospin orientations of the ground state nucleon ($\nu = p\uparrow, p\downarrow, n\uparrow, n\downarrow$). The wave function $\Psi_n(123)$ of the nucleon can be found as the ground state solution of the three-quark Hamiltonian

$$H(123) = \sum_{i=1}^{3} \left(m + \frac{p_i^2}{2m} \right) + \sum_{i<j=1}^{3} V_{ij} \qquad (6)$$

with the harmonic oscillator confinement potential (2). Since the Hamiltonian (6) is independent of spin, flavor and color (SFC) of the constituent quarks, the SFC-part $\chi_\nu(123)$ can be separated from the orbital part $\varphi_p(123)$ of the nucleon wavefunction according to

$$\Psi_n(123) = \varphi_p(123)\, \chi_\nu(123). \qquad (7)$$

The property of antisymmetry of the three-quark wavefunction determines the symmetry properties of the φ_P and the χ_ν part. In a systematic way, this decomposition can be done by using the technique of Young tableaux. The lowest energy eigenvalue corresponds to a total symmetric orbital part, with respect to spin and flavor the wave function has a mixed symmetry, whereas for the color part a total antisymmetric function is needed, see also [10]. With explicit notation of the spin ($\uparrow, \downarrow$), flavor (u, d) and color (R,G,B) degrees of freedom, the SFC-part of the nucleon wave function reads

$$\begin{aligned}\chi_\nu(123) = \frac{1}{\sqrt{18}} \{ & 2u\uparrow u\uparrow d\downarrow + 2u\uparrow d\downarrow u\uparrow + 2d\downarrow u\uparrow u\uparrow \\ & - u\uparrow u\downarrow d\uparrow - u\uparrow d\uparrow u\downarrow - d\uparrow u\uparrow u\downarrow - u\downarrow u\uparrow d\uparrow \\ & - u\downarrow d\uparrow u\uparrow - d\uparrow u\downarrow u\uparrow \} \frac{1}{\sqrt{6}} \det |RGB| . \end{aligned} \qquad (8)$$

By alternating the spin or isospin orientations in (8), the

four species of ground state nucleons ($\nu = p\uparrow, p\downarrow, n\uparrow, n\downarrow$) are described. The orbital part of the nucleonic wave function is obtained by solving the Schrödinger equation

$$H(123)\, \varphi_p(123) = E_n\, \varphi_p(123) \quad ; \quad n = P, \nu \, , \qquad (9)$$

yielding for the ground state

$$\varphi_p(123) = \frac{8\pi^3}{\Omega} \left(\frac{\sqrt{3}\, b^2}{\pi} \right)^{3/2} \delta_{P,P_R}\, e^{-\frac{b^2}{2}(p_\rho^2 + p_\lambda^2)} \, , \qquad (10)$$

$$E_n = \frac{P^2}{6m} + 3m + 3\sqrt{3}\,\omega \, . \qquad (11)$$

Here we have used the Jacobi coordinates

$$\begin{aligned} P_R &= p_1 + p_2 + p_3 \\ p_\rho &= \tfrac{1}{\sqrt{2}}(p_1 - p_2) \\ p_\lambda &= \tfrac{1}{\sqrt{6}}(p_1 + p_2 - 2p_3) \end{aligned} \, , \qquad (12)$$

and the width parameter of the gaussian wave function $b^{-2} = \sqrt{3}\, m\omega$; $\hbar = 1$.

Now, the antisymmetrized two-nucleon wave function follows from (5) with (7),(8) and (10). The normalisation is given by

$$\begin{aligned} N_{nn'} = \langle \Phi_{nn'} | \Phi_{nn'} \rangle = 1 - \delta_{P,P'} - 3 \sum_{p_1 \ldots p_6} & \varphi_P^*(123)\, \varphi_{P'}^*(456) \\ & * \{ c_{\nu\nu'}\, \varphi_p(126) \varphi_{p'}(453) - d_{\nu\nu'}\, \varphi_p(453) \varphi_{p'}(126) \} . \end{aligned} \qquad (13)$$

Here, the $c_{\nu\nu'}$ and $d_{\nu\nu'}$ reflect the scalar products of the SFC-part with exchange according to $P_{3,6}$ and $P_{nn'}P_{3,6}$

$$\begin{aligned} c_{\nu\nu'} &= \langle \chi_\nu(123)\, \chi_{\nu'}(456)\, \chi_\nu(126)\, \chi_{\nu'}(453) \rangle \\ d_{\nu\nu'} &= \langle \chi_\nu(123)\, \chi_{\nu'}(456)\, \chi_\nu(453)\, \chi_{\nu'}(126) \rangle \, . \end{aligned} \qquad (14)$$

The color degrees of freedom are immediately elaborated by rearranging the color variables, a factor 2 arises from two different variants of χ_ν if the non exchanged variables are transposed. The remaining SF-variables are explicitely written down and evaluated. The results for $\nu = p\uparrow$ are given in the Table, the equivalent results hold also for the other nucleon states, if the interaction is invariant with respect to the isospin variables.

The momentum variables are integrated taking into account that the exchange operator $P_{3,6}$ is different from zero only for $p_3 = p_6$. The result can be given in a closed form

$$N_{nn'} = 1 - \delta_{\ell,\ell'} - \frac{9\sqrt{3}}{8}\left(\frac{b^2}{\pi}\right)^{3/2}\frac{\delta\pi^3}{\Omega}\left\{c_{\nu\nu'}\, e^{-\frac{b^2}{12}(P-P')^2} - d_{\nu\nu'}\, e^{-\frac{b^2}{8}(P-P')^2}\right\} \qquad (15)$$

Table 1. The values of the matrix elements $c_{\nu\nu'}$ and $d_{\nu\nu'}$ for $\nu = p\uparrow$.

ν	ν'	$c_{\nu\nu'}$	$d_{\nu\nu'}$
$p\uparrow$	$p\uparrow$	31/243	31/243
	$p\downarrow$	14/243	17/243
	$n\uparrow$	14/243	17/243
	$n\downarrow$	22/243	25/243
	$\sum_{\nu'}$	1/3	10/27

3 Evaluation of the Energy Shift

The antisymmetrisation of the two-nucleon wave function with respect to the quark degrees of freedom leads to a shift in the two-nucleon energy according to

$$\Delta E^{\text{Pauli}}_{nn'} = \frac{1}{N_{nn'}}\langle\phi_{nn'}|H|\phi_{nn'}\rangle - E_n - E_{n'} \,, \qquad (16)$$

with E_n given by eq. (11). The Hamiltonian H = KE + PE contains the kinetic part, eq. (1), and the potential part, eq. (3). Neglecting the antisymmetrisation of the wave function with respect to quark exchange, the kinetic energy which is in Jacobi coordinates

$$KE = 6m + \frac{p_\varrho^2}{6m} + \frac{p_\varrho'^2}{6m} + \frac{1}{2m}\left(p_\varsigma^2 + p_\lambda^2 + p_\varsigma'^2 + p_\lambda'^2\right) \,, \qquad (17)$$

and the potential energy

$$PE = \frac{m\omega^2}{2}\cdot 3\left(\varsigma^2 + \lambda^2 + \varsigma'^2 + \lambda'^2\right) \,, \qquad (18)$$

are immediately evaluated for the two-nucleon system with the result

$$\langle\phi_{nn'}|KE|\phi_{nn'}\rangle \approx 6m + \frac{p^2}{6m} + \frac{p'^2}{6m} + 3\sqrt{3}\,\omega \,, \qquad (19)$$

$$\langle\phi_{nn'}|PE|\phi_{nn'}\rangle \approx 3\sqrt{3}\,\omega \,, \qquad (20)$$

so that no energy shift arises.

Orthogonalisation of the wave function by antisymmetrisation will lead to a change in the kinetic energy. In contrast to the kinetic energy, the potential energy (3) is not determined by the wave function but by the density distribution of the quarks. In particular, the probability of a given quark configuration is determined by the density distribution function. It is well-known from the Hartree-Fock theory that antisymmetrisation will not change the particle density distribution $\varrho(\underline{r}) = \sum_i \delta(\underline{r}-\underline{r}_i)$. For two nucleons we obtain an overlap of the quark density distributions, and the potential energy in not significantly changed by the antisymmetrisation procedure, as long as string flip processes are not of importance. As discussed below, a variation of the wave function beyond the scope of a Hartree-Fock type antisymmetrisation will also lead to a variation of the potential energy.

In this way, the energy shift (16) is determined by the change of the kinetic energy

$$\Delta E^{Pauli}_{nn'} = N_{nn'}^{-1} \{ E_n + E_{n'} + \Delta KE^{Pauli}_{nn'} \} - E_n - E_{n'} , \quad (21)$$

with

$$\Delta KE^{Pauli}_{nn'} = - (E_n + E_{n'}) \delta_{p,p'} - 3 \sum_{p_1 \cdots p_6} \varphi^*_p(123)\, \varphi^*_{p'}(456) \quad (22)$$
$$* \; KE \{ c_{\nu\nu'} \varphi_p(126) \varphi_{p'}(453) - d_{\nu\nu'} \varphi_p(453) \varphi_{p'}(126) \} .$$

Expanding the normalisation factor $N_{nn'}^{-1}$ up to the first order with respect to the overlap integral, see eq. (13), the expression (4) for the Pauli shift is recovered.

The interpretation of the energy shift due to the Pauli blocking can be given in correspondence to atomic physics. At short interatomic distances, the energy of the two-atom system is sharply increasing what is usually represented by a repulsive, hard core like interaction potential. Indeed, the physical reason of this increase of energy is not the Coulombic electron-electron interaction, but the increase of kinetic energy because of the Pauli principle which demands the orthogonalisation of the electron wave functions.

Now, let us proceed to the explicit evaluation of the Pauli-blocking shift (21) which may be given using the Jacobi coordinates (12) as follows

$$\Delta E^{Pauli}_{nn'} = 3 \sum_{P_R P_S P_\lambda} \sum_{P'_R P'_S P'_\lambda} \frac{\partial(p_1 \cdots p_6)}{\partial(P_R \cdots P'_\lambda)} \delta_{P,P_R} \delta_{P',P'_R} \{ 6\sqrt{3}\,\omega$$
$$- \frac{1}{2m} (p_S^2 + p_\lambda^2 + p_S'^2 + p_\lambda'^2) \} \{ \delta_{P,P_R} - \tfrac{1}{3}(P_R - P'_R) + \tfrac{2}{\sqrt{6}}(p_\lambda - p'_\lambda)$$
$$- \delta_{P,P'_R} + \tfrac{1}{3}(P_R - P'_R) - \tfrac{2}{\sqrt{6}}(p_\lambda - p'_\lambda) \} \exp\{ - b^2 (p_S^2 + p_\lambda^2 + p_S'^2 + p_\lambda'^2) \}$$
$$= \frac{9\sqrt{3}}{16} \frac{8\pi^3}{\Omega} \left(\frac{b^2}{\pi}\right)^{3/2} \frac{1}{mb^2} \{ c_{\nu\nu'} e^{-\frac{b^2}{12}(P-P')^2} [\tfrac{15}{2} - \tfrac{b^2}{12}(P-P')^2]$$
$$- d_{\nu\nu'} e^{-\frac{b^2}{3}(P-P')^2} [\tfrac{15}{2} - \tfrac{b^2}{3}(P-P')^2] \} . \quad (23)$$

Whereas this quantity (23) measures the surplus energy arising from the antisymmetrisation of the wave function with respect to the two-nucleon problem, we are especially interested in the energy shift for a single nucleon ΔE^{Pauli}_n in a many-nucleon system which can be obtained from (23) by summation over the second nucleonic index n' whereby at T=0 the respective distribution function (see (4)) is a step function restricting the momentum summation to the range of the Fermi sphere $|P'| < P_F$. The sum over P' may then be evaluated as an integral yielding

$$\Delta E^{Pauli}_n = \frac{\Omega}{8\pi^3} \int_{|P'|<P_F} d^3P' \sum_{\nu'} \Delta E^{Pauli}_{nn'}$$
$$= \frac{3\sqrt{3}}{16\sqrt{\pi}} \frac{b}{m} \int_0^{P_F} dP' P'^2 \int_{-1}^{1} dz \{ e^{-\frac{b^2}{12}(P^2+P'^2-2PP'z)} [\tfrac{15}{2} - \tfrac{b^2}{12}(P^2+P'^2-2PP'z)]$$
$$- \frac{10}{9} e^{-\frac{b^2}{3}(P^2+P'^2-2PP'z)} [\tfrac{15}{2} - \tfrac{b^2}{3}(P^2+P'^2-2PP'z)] \}$$
$$= \frac{3\sqrt{3}}{32\sqrt{\pi}} \frac{1}{mbP} \int_{-P_F}^{P_F} dP' P' \{ e^{-\frac{b^2}{12}(P-P')^2} [156 - 26b^2(P-P')^2]$$
$$- \frac{10}{9} e^{-\frac{b^2}{3}(P-P')^2} [39 - 26b^2(P-P')^2] \} . \quad (24)$$

In the low density limit, i.e. for small values of P_F, the integrand of eq.(24) can be expanded with respect to P,P'. Performing the P'- integral, for the Pauli shift of a nucleon with the momentum P follows

$$\Delta E^{Pauli}_{\nu P} = \frac{5}{8\sqrt{3\pi}} \frac{b}{m} \left\{ - P_F^3 + \frac{527}{60} b^2 \left(\frac{1}{5} P_F^5 + \frac{1}{3} P_F^3 P^2 \right) \right\} . \quad (25)$$

Especially, at the Fermi momentum P_F we have

$$\Delta E^{Pauli}_{\nu P_F} = \frac{5}{8\sqrt{3\pi}} \frac{b}{m} \left\{ - P_F^3 + \frac{1054}{225} b^2 P_F^5 \right\} . \quad (26)$$

This shift is at T=0 identical with the shift in the chemical

potential and can be used to derive the equation of state , see Ref.[5].

In order to give numerical results for the Pauli shift (24), we adopt the values m=350 MeV and b=0.59 fm according to Ref. 3 which reproduce quite well the single nucleon properties. With the relation $P_F^3 = (3\pi^2/2)\varrho$, the Pauli blocking shift can be given as a function of the nuclear matter density ϱ

$$\Delta E^{Pauli}(\varrho) = a_1 \varrho + a_2 \varrho^{5/3} \quad , \qquad (27)$$

with $a_1 = -197.77$ MeV fm^3 and $a_2 = 1944.45$ MeV fm^5. This expression (27) is presented in Fig.1.

4. Consequences for the Properties of Nuoleons in Nuclear Matter

The energy shift of a nucleon in dense nuclear matter due to the Pauli quenching on the quark substructure can be related to other effects: the density dependence of the energy of nuclear matter, the effective nucleon mass, the effective nuoleon-nuoleon interaction potential, and the swelling of nuclear radii.

We first discuss the density dependence of the nuclear matter energy at zero temperature, for which a phenomenological parametrisation is given by the Skyrme interaction [11]

$$\Delta E^{Skyrme}(\varrho) = b_1 \varrho + b_2 \varrho^{5/3} + b_3 \varrho^2 \quad , \qquad (28)$$

with $b_1 = -792.97$ MeV fm^3, $b_2 = 125.225$ MeV fm^5, $b_3 = 2711.9$ MeV fm^6. This parametrisation reproduces the binding property of nuclear forces represented by the negative value of b_1, and a hard core contribution which leads to the saturation at nuclear matter density.

Of course, the Pauli quenching shift cannot reproduce the binding part of the Skyrme interaction because no virtual meson exchange

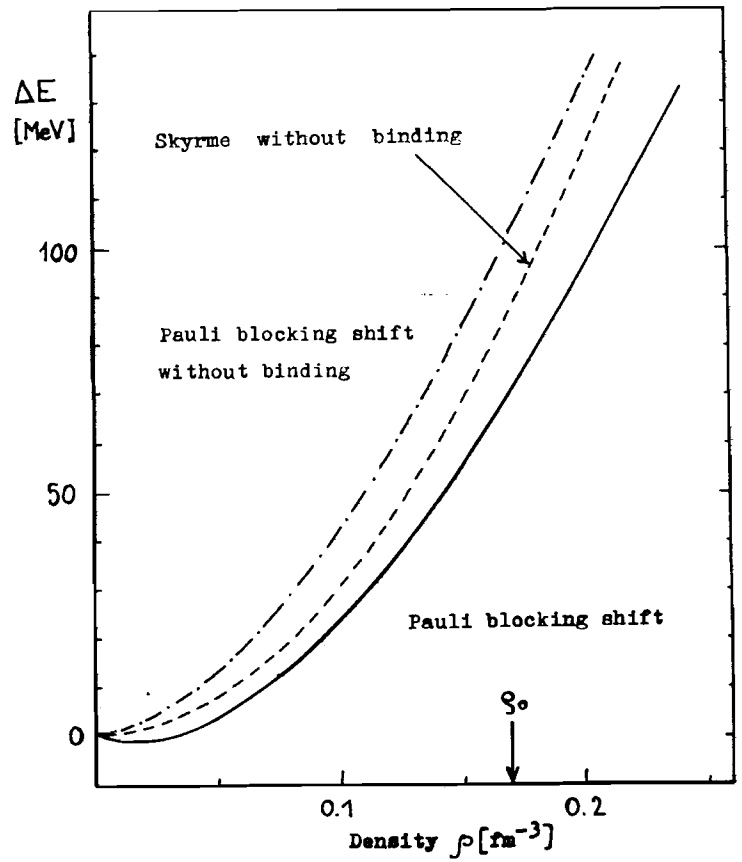

Fig. 1. The Pauli blocking shift ΔE^{Pauli} for a nucleon at T =0 in dependence on the density ϱ of the surrounding nuclear medium (full line). For comparison, the Skyrme parametrisation without binding (dashed line) and the Pauli blocking shift without binding (dot-dashed line) are also shown.

is considered within this approach. For this, one has to take into account besides the six-quark configurations for two nucleons also configurations with additional quark-antiquark pairs. Therefore it makes sense to compare the energy shifts without the binding part

only, i.e. without the term linear in density. As is shown in Fig.1, a fairly good agreement is observed so that we conclude that the strong increase of the nuclear matter energy at high densities which is usually interpreted by the hard-core like repulsion between nucleons at short distances can be related to the Pauli blocking effect on the quark substructure level. The analogous effect is well known from atomic physics where a pseudopotential for the electron-ion interaction is constructed so that the Pauli principle leads to a hard-core like effective repulsive potential.

Within this quark-substructure derivation of the nuclear matter energy shift, it would also be possible to investigate the effect of finite temperatures. On the other side, the results given here should be considered as a first model calculation to be improved by the use of more realistic quark potentials and the account of virtual quark pair excitations.

The effective mass of nucleons in nuclear matter can be derived from the Pauli quenching shift (4) which can be interpreted as the self-energy of the nucleon. The temperature dependence of the effective mass and the comparison to other approaches has been discussed in Ref. [6].

The consequences of the Pauli principle on the quark-substructure level are also of interest for the behaviour of scattering phases at high energies. We will not discuss the relation to effective short-range potentials as, e.g., a hard-core repulsive potential in this paper. The effects of antisymmetrisation on the effective nucleon-nucleon potential constructed from a treatment of the six-quark problem were extensively studied by several authors, see Refs. [7-9]. Within a few-particle treatment the Resonating group method has been applied. In general, hidden color states are considered which are excluded in the present work. Using realistic quark inter-

action potentials, quite good results are obtained for the lowest nucleon-nucleon scattering phase shifts [7-9].

Another interesting effect of the quark substructure of nucleons is the swelling of nucleonic radii as discussed in connection with the EMC effect [12]. The relation to the Pauli quenching effect was discussed in [13] and will be given here starting from the results of this paper.

From (27) we find the energy shift per nucleon in nuclear matter at zero temperature and density

$$\Delta E = \frac{1}{\varrho} \int_0^{\varrho} \Delta E^{Pauli}(\varrho')\, d\varrho'$$

$$= \frac{1}{2} a_1 \varrho + \frac{3}{8} a_2 \varrho^{5/3} \quad , \qquad (29)$$

which indicates that the energy of a single nucleon increases with the density of matter.

Since we are dealing with a quadratic interaction the potential energy is related to the mean square radius of the nucleon (mean string length). According to the virial theorem the total energy is divided into equal parts for the kinetic and potential energies, respectively. However, the mean values have to be evaluated with the actual wave functions affected by the surrounding medium. Therefore, the relative change of total energy coincides with the relative change of the potential energy. Denoting by E_0 the energy of a nucleon at rest and at zero density, one has the relation

$$\langle r^2 \rangle / \langle r^2 \rangle_0 = E / E_0 = 1 + \Delta E / E_0 \quad , \qquad (30)$$

where $\langle r^2 \rangle_0$ refers to the free nucleon. With the parameter values given above the change in the effective radius of the nucleon is described by the relation

$$\langle r^2\rangle / \langle r^2\rangle_0 = 1 + mb^2 \left(\frac{a_1}{6}\rho + \frac{a_2}{8}\rho^{5/3} \right) , \qquad (31)$$

and is shown in Fig. 2. Especially, the increase of the average radius of a nucleon immersed in nuclear matter of the density $\rho_0 = 0.17\ \mathrm{fm}^{-3}$ amounts to

$$\langle r^2\rangle / \langle r^2\rangle_0 = 1.022 \qquad (32)$$

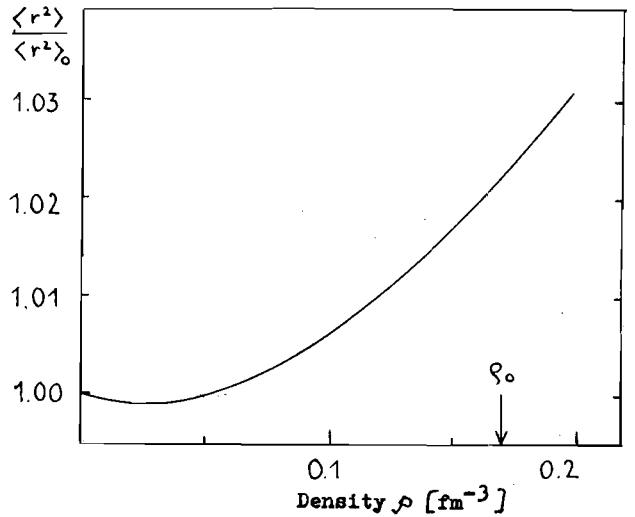

Fig. 2: The density dependence of the effective nucleon radius in nuclear matter at T = 0. At nuclear matter density $\rho_0 = 0.17\ \mathrm{fm}^{-3}$ a swelling of about 2% is obtained.

This enlargement of the average nucleon radius agrees fairly well with that advocated in Refs. [9,12]. The results illustrate that the Pauli blocking mechanism becomes more operative with increasing medium density. An increase of the swelling of nucleon radii with increasing temperature was obtained in Ref.[13] for different densities.

Of course, the results shown in Fig. 2 should be considered only as an estimation. For the quantitative comperison with experiments a realistic quark interaction potential should be used instead of the harmonic oscillator potential and the actual form of the wave function should be considered instead of using the virial theorem. The consequence of the antisymmetrisation of the quark wave function has been recently considered [14] in order to explain the EMC-effect. However, we will not discuss the many other approaches to find a theoretical explanation of the EMC-effect.

In conclusion, an important effect of the quark substructure is the Pauli quenching effect in dense matter. Within a many-quark treatment an energy shift of nucleons is obtained from the antisymmetrisation procedure for the many-quark wave function. Typical effects of the surrounding dense matter are the density dependent energy (equation of state), the effective mass of nucleons and the swelling of nucleonic radii. Within our quantum statistical approach, these effects can also be considered at finite temperatures.

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Паули-квенчинг адронов в ядерной материи -
эффект кварковой субструктуры

На основе концепции насыщения струнного взаимодействия между кварками в многокварковой системе рассматривается эффект запрета Паули в адронной ядерной материи. "Паули-квенчинг" связанных состояний приводит к сдвигу энергии. Результат сравнивается со случаем жесткого кора потенциала Скирма. Для связанных состояний кварков вычисляется возрастание эффективного радиуса нуклона в ядерной материи за счет запрета Паули.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1988

Blaschke D., Röpke G. E2-88-77
Pauli-Quenching for Hadrons in Nuclear Matter -
a Quark Substructure Effect

Starting from a string saturation model of quark interaction in a many-quark system, the Pauli quenching in the bound state phase is considered which leads to an energy shift comparable with the hard-core part of the Skyrme interaction. The Pauli quenching of quark bound states in nuclear matter can be related to the increase of the effective radii of nucleons.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Preprint of the Joint Institute for Nuclear Research. Dubna 1988