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Short Range Part of Baryon-Baryon Interaction in a Quark Model. II

—— *Numerical Results for S-Wave* ——

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An approach to the short range part of baryon-baryon interaction based on a non-relativistic quark model, proposed and formulated in the previous paper, is applied to the case of S-wave relative motion. Repulsive core like interactions appear in most of the spin(S)-isospin(T) states, including the NN system with $(S, T)=(1, 0)$ and $(0, 1)$, while attractive interactions are predicted for the $\Delta\Delta$ system with $(S, T)=(3, 0)$ and $(1, 0)$. The interaction is mostly due to the Pauli principle between quarks and the spin-spin term in the quark-quark force, and is insensitive to the confinement term. The effect of meson exchanges, which is not included in the present non-relativistic quark model, is studied by simulating it with a phenomenological local potential. Qualitative behaviours of the S-wave NN phase shifts are reproduced in this way.

§ 1. Introduction

In the previous paper,[1] hereafter referred to as (I), we proposed to explain the short range part of the baryon-baryon interaction, on the basis of non-relativistic quark model. The bound state and scattering problems of two baryon systems were formulated by the resonating group method (RGM),[2] which had been extensively developed in the nuclear cluster model calculations. A baryon is considered as a three quark cluster and is described by a wave function of the form

$$\phi_A(\xi_A)=C([111]:\xi_A{}^C)\mathscr{S}([3]ST:\xi_A{}^S)\ \varphi_A(\boldsymbol{\xi}_1, \boldsymbol{\xi}_2), \tag{1·1}$$

where C is the color part and is assumed to be a totally antisymmetric [111] color singlet state, while $\mathcal{S}$ is the spin-isospin part and is given by a totally symmetric [3] state with $S=T=1/2$ for nucleon (N) and $S=T=3/2$ for isobar (Δ). The orbital part φ_A is taken to be of the Gaussian form,

$$\varphi_A(\boldsymbol{\xi}_1, \boldsymbol{\xi}_2)=\left(\frac{1}{2\pi b^2}\right)^{3/4}\left(\frac{1}{3\pi b^2}\right)^{3/4}\exp\left\{-\left(\frac{\xi_1{}^2}{4b^2}+\frac{\xi_2{}^2}{3b^2}\right)\right\}, \tag{1·2}$$

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where $\boldsymbol{\xi}_1$ and $\boldsymbol{\xi}_2$ are the internal coordinates of the three quarks defined by (see Eq. (4·1) of (I))

$$\boldsymbol{\xi}_1 \equiv \boldsymbol{r}_1 - \boldsymbol{r}_2, \ \boldsymbol{\xi}_2 \equiv \boldsymbol{r}_3 - \frac{\boldsymbol{r}_1 + \boldsymbol{r}_2}{2} . \tag{1·3}$$

The RGM assumes that a system of two baryons, A and B, is described by a wave function of the form

$$\Psi(\boldsymbol{\xi}_A, \boldsymbol{\xi}_B, \boldsymbol{R}_{AB}) = \mathcal{A}[\phi_A(\boldsymbol{\xi}_A)\phi_B(\boldsymbol{\xi}_B)\chi(\boldsymbol{R}_{AB})], \tag{1·4}$$

where χ is a wave function for the relative motion and $\mathcal{A}$ is the anti-symmetrization operator for the six quarks.

The hamiltonian for the quark system[4),5)] consists of the kinetic energy, K, and the interaction, V, i. e.,

$$H = K + V , \tag{1·5}$$

where

$$K = \Sigma \frac{p_i^2}{2m_q} - K_G ,$$

$$V = V^{\mathrm{CONF}} + V^{\mathrm{OGEP}} \tag{1·6}$$

with

$$V^{\mathrm{CONF}} = \sum_{i<j} (-a)(\lambda_i \cdot \lambda_j) r_{ij}$$

or

$$= \sum_{i<j} (-a')(\lambda_i \cdot \lambda_j) r_{ij}^2 \tag{1·7}$$

and

$$V^{\mathrm{OGEP}} = \sum_{i<j} (\lambda_i \cdot \lambda_j) \frac{\alpha_s}{4 r_{ij}} - \frac{\pi \alpha_s}{6 m_q^2} (\lambda_i \cdot \lambda_j)(\boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j) \boldsymbol{\delta}(\boldsymbol{r}_{ij}). \tag{1·8}$$

The RGM equation for the relative wave function χ is determined by

$$\int \phi_A^\dagger(\boldsymbol{\xi}_A)\phi_B^\dagger(\boldsymbol{\xi}_B)(H-E)\, \Psi(\boldsymbol{\xi}_A, \boldsymbol{\xi}_B, \boldsymbol{R}_{AB}) d\boldsymbol{\xi}_A d\boldsymbol{\xi}_B = 0$$

which gives

$$\int [H(\boldsymbol{R}, \boldsymbol{R}') - EN(\boldsymbol{R}, \boldsymbol{R}')]\chi(\boldsymbol{R}') d\boldsymbol{R}' , \tag{1·9}$$

where $H(\boldsymbol{R}, \boldsymbol{R}')$ and $N(\boldsymbol{R}, \boldsymbol{R}')$ are the hamiltonian and normalization kernels, respectively, and are defined in (I) (Eqs. (4·5)~(4·12)). General discussions on the properties of these kernels are also given in (I). A variational method[3)] was used to solve this equation.

For an illustrative purpose, a renormalized wave function χ_R is introduced by

$$\chi_R \equiv \int N^{1/2}(\boldsymbol{R}, \boldsymbol{R}')\chi(\boldsymbol{R}')d\boldsymbol{R}' \tag{1·10}$$

and the equivalent local potential $V^{EQ}(R)$ is defined as follows:

$$\left[-\frac{1}{2\mu}\frac{d^2}{dR^2}+\frac{L(L+1)}{2\mu R^2}+V^{EQ}(R)\right]\chi_R{}^L(R)=E_{\rm rel}\chi_R{}^L(R). \tag{1·11}$$

In this paper, we report some results of the calculation based on the formalism of (I). We first discuss how to determine the several parameters of the model (§ 2). Numerical results of the scattering and the bound state problems for the S-wave two baryon states are given in § 3. Effects of the meson exchange contributions in the intermediate and long range part are also investigated in § 3. Discussion and conclusion are given in § 4.

§ 2. Determination of parameters

The harmonic oscillator constant b can be determined by the root-mean-square (RMS) charge radius of a proton. The RMS radius for the $(0s)^3$ configuration of quarks is given by

$$\sqrt{\langle r^2\rangle}=\sqrt{\frac{3}{2}\left(1-\frac{1}{3}\right)b^2}=b\,, \tag{2·1}$$

which gives $b=0.8$ fm. We, however, consider a nucleon to have meson cloud around the core of three quarks, so that the actual extension of the "quark core" part may be smaller than the above value. In the present calculation, we set $b=0.6$ fm as a standard value and discuss the effect of changing b around this value. Further, we assume that an isobar Δ has the same extension as a nucleon, mainly because of simplicity.

Masses of N and Δ are obtained as the expectation values of the hamiltonian for the $(0s)^3$ configuration of quarks. With the quark rest mass $3m_q$, they are given by

$$M_N=3m_q+\frac{3}{2m_qb^2}-8\varepsilon-\frac{2}{2\sqrt{2\pi}}\frac{\alpha_s}{m_q{}^2b^3}\,, \tag{2·2a}$$

$$M_\Delta=3m_q+\frac{3}{2m_qb^2}-8\varepsilon+\frac{2}{3\sqrt{2\pi}}\frac{\alpha_s}{m_q{}^2b^3}\,, \tag{2·2b}$$

where

$$\varepsilon = \frac{\alpha_s}{2\sqrt{2\pi}\,b} - \begin{cases} 2\sqrt{\dfrac{2}{\pi}}\,ab & \text{for the linear confinement,} \\ 3a'b^2 & \text{for the quadratic confinement.} \end{cases} \tag{2·3}$$

We next relate the strength of the confinement potential a or a' to the extension parameter b, regarding V^{CONF} as the main part of the potential. For the quadratic confinement potential, our wave function Eq. (1·2) is an eigenfunction of the hamiltonian without V^{OGEP} if

$$a' = \frac{1}{16 m_q b^4} \,. \tag{2·4}$$

For the linear one, we will determine a by making the baryon mass M_B without V^{OGEP} stationary at a given value of b. Such a variational consideration gives

$$a = \frac{3\sqrt{2\pi}}{32} \frac{1}{m_q b^3} \,. \tag{2·5}$$

The relations (2·4) and (2·5) may well be modified due to the perturbation V^{OGEP}. This modification, however, does not give rise to any serious problem, since the results of numerical calculations are almost independent of a or a', as we will see later.

The quark-gluon coupling constant α_s can be determined in several ways. Though α_s depends on the energy of applied system or the distance between two quarks according to QCD,[4] we will neglect such effects mainly because the dependence is logarithmic and will not be so serious. In the present calculation, α_s is determined from the mass difference between N and Δ as

$$M_\Delta - M_N = \frac{4}{3\sqrt{2\pi}} \frac{\alpha_s}{m_q{}^2 b^3} \,, \tag{2·6}$$

since we want to have the correct threshold difference between NN, $N\Delta$ and $\Delta\Delta$ channels in the scattering problems. The mass difference in the present model is entirely due to the spin-spin (color magnetic) part of V^{OGEP} in Eq. (1·8), but the contribution of the meson cloud may again modify the relation (2·6). With this in mind, we will see the effect of changing α_s for the Coulomb (spin independent) part of V^{OGEP}.

In Table I, actual values of the parameters determined by Eqs. (2·4)~(2·6) are tabulated for $b=0.6$, 0.8 and 0.5 fm. The last two columns of Table I show the corresponding masses of N and Δ given by Eqs. (2·2a) and (2·2b), respectively. We have no intention of reproducing the observed baryon masses in the present model, since, first, earlier works[5] show that one needs an additional constant term in the hamiltonian to fit the hadron masses, second, meson cloud may give appreciable contribution and finally we feel that the nonrelativistic model will not be sufficient to discuss the total energy of the system.

Table I. Six sets of parameters with corresponding mass of N and Δ. The first set (I) is mainly adopted in the present calculation.

	b(fm)	α_s	a(MeV/fm)	a'(MeV/fm^2)	M_N(MeV)	M_Δ(MeV)
I	0.6	1.39	141.	—	1645	1938
I′	0.6	1.39	—	62.5	1102	1396
II	0.8	3.30	59.6	—	368	661
II′	0.8	3.30	—	19.8	61	335
III	0.5	0.81	244.	—	4080	4372
III′	0.5	0.81	—	129.8	2862	3155

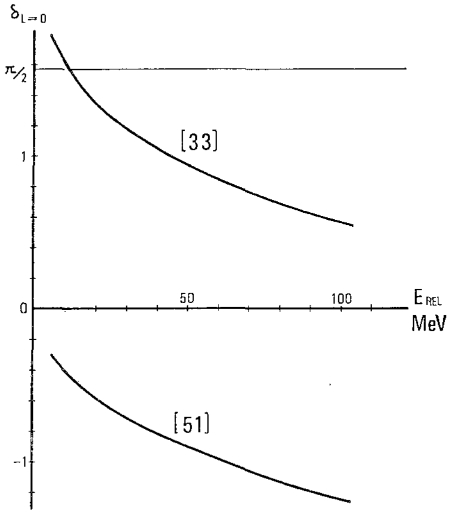
It should also be noted here that the masses for baryons given above have nothing to do with the mass in the kinetic energy of the relative motion between two baryons which is always $\frac{3}{2}m_q$ in the present approach. Only the mass difference between N and Δ, which determines threshold energy differences for NN, $N\Delta$ and $\Delta\Delta$ channels, is relevant in the following calculations.

A comment should be given on the one-gluon exchange potential V^{OGEP} used in this paper. The values of α_s determined by the relation (2·6) are somewhat larger than those usually expected in QCD at the squared momentum transfer $Q^2 \sim 1(\text{GeV})^2$, namely $\alpha_s(Q^2=1)\sim 1$. It is also hard to justify that one can use a perturbative approach for the quark-quark interaction with such a strong "hyperfine" constant α_s. On the other hand, the comparison between the results of Refs. 6) and 7) and the corresponding ones in the present calculation indicates that the specific form of the quark-quark interaction is almost irrelevant. It is therefore more reasonable to consider V^{OGEP} used here as an effective interaction including higher order effects.

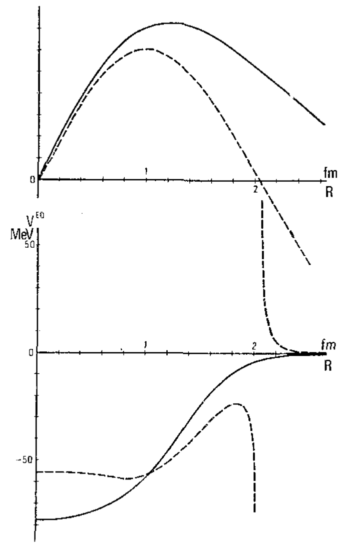
§ 3. Numerical results

3.1. *SU*(4) *invariant limit*

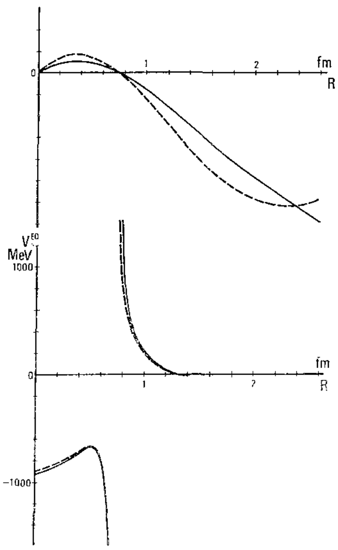
We consider first a simple case where the spin-spin term is omitted from the hamiltonian while the "Coulomb" term is included. In this case, the hamiltonian is invariant under any transformation of the group $SU(4)$ and the masses of N and Δ are completely degenerate. We have only two different $SU(4)$ channels specified by the symmetry, i. e., [33] and [51] eigenchannels. Figures 1(a)~(c) show the calculated phase shifts, the renormalized wave functions, χ_R's, and the equivalent local potentials (EQLP) for these cases. The phase shift and the EQLP for the [33] channel indicate the existence of a bound state in this channel, which is indeed found at a binding energy $E_B=1.9$ MeV. On the contrary, the phase shift for the [51] channel shows that the interaction is repulsive in this channel. The renormalized wave function has an almost energy independent node at about $R=0.74$ fm, which reflects the existence of a forbidden $0s$ state to which χ_R must be orthogonal. This feature suggests that the effective interaction between two baryons has a strongly repulsive core at the node, which is more clearly seen in the behavior of the EQLP. Further, one can easily find that the phase shift is similar to that of the scattering from a hard sphere, i. e., $\delta\sim -kd$. Here d is the core radius and corresponds to the position of the node of the wave



(a)



(b)



(c)

Fig. 1.(a) Baryon-baryon S wave phase shifts for the [33] and [51] symmetric $SU(4)$ eigenstates when the spin-spin interaction is omitted. Parameters are those of the set I($b=0.6$ fm, linear confinement) in Table I.

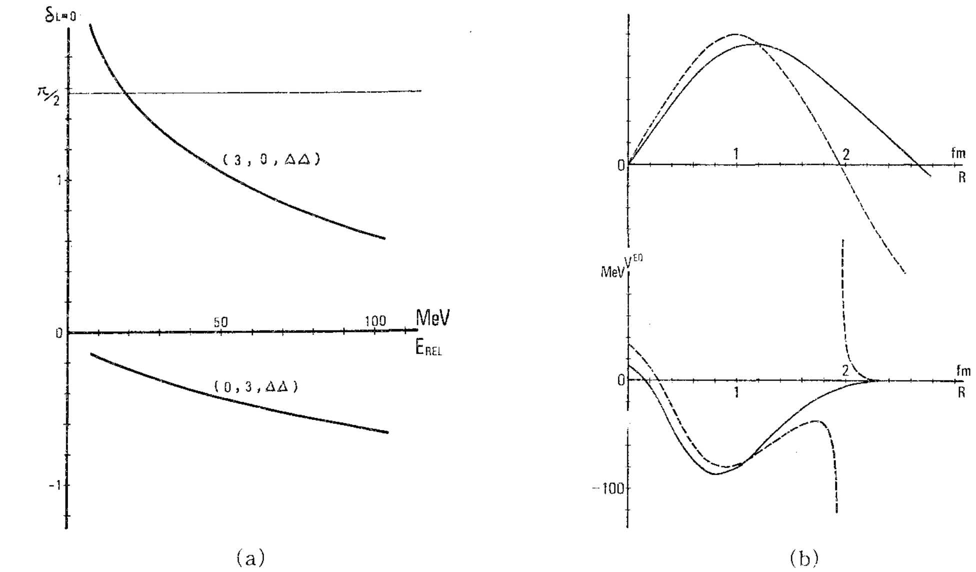
(b) Wave functions $\chi_R(R)$ (the upper part) and equivalent local potentials (EQLP) $V^{EQ}(R)$ (the lower part) at $E_{rel}=10$ MeV (full line) and 50 MeV (dashed line) for the [33] symmetric state.

(c) Wave functions and EQLP's for the [51] symmetric state.

function in this case. These features are to be expected from the existence of the forbidden state, as a similar situation occurs in the RGM calculation of, for instance, the α-α scattering.[8]

3.2. [33] *eigenstates*

When the spin-spin interaction is turned on, the hamiltonian becomes non-invariant with respect to the $SU(4)$ transformation and the masses of N and Δ split. The degeneracies of the states in Table II of (I) also disappear. Among them the states $(3, 0, \Delta\Delta)$ and $(0, 3, \Delta\Delta)$ still belong to the [33] representation.



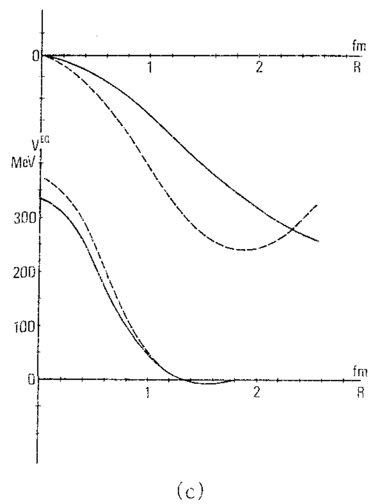


Fig. 2. (a) Phase shifts for the $(3, 0, \Delta\Delta)$ and $(0, 3, \Delta\Delta)$ states.

(b) Wave functions and EQLP's for the $(3, 0, \Delta\Delta)$ state.

(c) Wave functions and EQLP's for the $(0, 3, \Delta\Delta)$ state.

Figures 2(a)~(c) show the results of calculation for these states. The $(3, 0, \Delta\Delta)$ state has a bound state at $E_B = 6.6$ MeV for the parameter set I, while the bound state disappears and the interaction becomes repulsive for the $(0, 3, \Delta\Delta)$ state, although the corresponding wave function has no energy independent node. These results indicate that the spin-spin interaction plays an important role in the two baryon problems.

The existence of a $\Delta\Delta$ bound state in the $(S, T)=(3, 0)$ channels is very interesting because Kamae et al. suggested a dibaryon resonance at $\sqrt{s}=2350$ MeV possibly in this channel.[9] The position of the resonance corresponds to the binding energy $E_B = 114$ MeV, which is much deeper than the prediction of the present model. One should, however, recall that the effect of the meson cloud has not been included in the calculation and the model predicts only the hard core like repulsion in most of the channels.

3.3. [51] *eigenstates*

The four states $(3, 2, \Delta\Delta)$, $(2, 3, \Delta\Delta)$, $(1, 1, N\Delta_a)$ and $(2, 2, N\Delta_a)$ have the unique [51] symmetric $SU(4)$ wave functions. The spin-spin interaction give little influence on these states. They have similar phase shifts, the wave functions and the EQLP's. The energy independent nodes of the relative wave functions, χ_R's, appear at

	$R=0.75$ fm	for	$(3, 2, \Delta\Delta)$,
	0.78 fm	for	$(2, 3, \Delta\Delta)$,
	0.74 fm	for	$(1, 1, N\Delta_a)$
and	0.77 fm	for	$(2, 2, N\Delta_a)$,

and correspond to the radii of the repulsive cores. None of them have a bound state. Figure 3 shows the phase shifts calculated for these states.

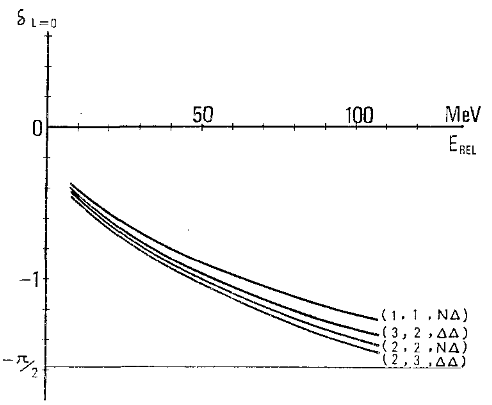
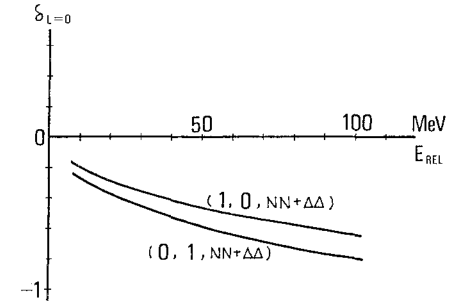


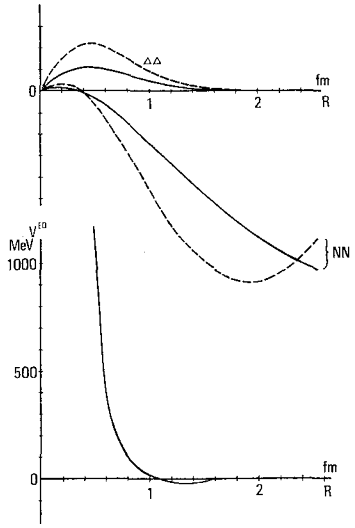
Fig. 3. Phase shifts for the states which belong to the [51] symmetric $SU(4)$ representation.

3.4. *Symmetry mixed states*

For the most interesting cases of $(S, T)=(1, 0)$ and $(0, 1)$, the [33] and [51] $SU(4)$ symmetries mix with each other and the coupled channel calculations are carried out. Figures 4(a)-(c) show the results, where the energy $E_{\rm rel}$ is measured from the two nucleon threshold. The behavior of the phase shifts and the EQLP's for the NN channel indicates the existence of a strong short range repulsion between two nucleons in both channles. The slope of the phase shifts $(-d\delta/dk)$



(a)



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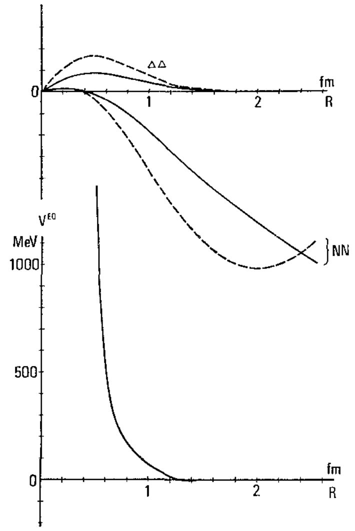


(c)

Fig. 4.(a) Phase shifts for the coupled (1, 0, $NN+\Delta\Delta$) and (0, 1, $NN+\Delta\Delta$) states.

(b) Wave functions and EQLP's of the N-N channel for the (1, 0, $NN+\Delta\Delta$) state.

(c) Wave functions and EQLP's of the N-N channel for the (0, 1, $NN+\Delta\Delta$) state.

averaged over the range 100 MeV $\leq E_{\rm rel} \leq$ 200 MeV is 0.39(0.45)fm for $^3S(^1S)$ state and corresponds to the radius of the repulsive core. The intermediate range attraction does not show up in the present calculation.

It should be noted that the origin of the repulsive cores obtained here is not the same as that in the [51] symmetric channels because the relative $0s$ state is allowed through the [33] symmetric $SU(4)$ state in the case of the mixed symmetry. In order to make it clear, calculations are made without coupling to the $\Delta\Delta$ channel. We obtain almost the same phase shifts except near the $\Delta\Delta$ threshold (Fig. 5(a)), while the wave functions and the EQLP's are rather different in the internal region (Figs. 5(b) and (c)). It indicates that the repulsive core is caused by a mechanism other than the symmetry structure of the wave function. On the other hand, if we omit the spin-sipn term in the quark-quark potential (1·8), the resulting phase shifts become very flat and the corresponding EQLP's become very weak. Thus, we conclude that the spin-spin (color magne-

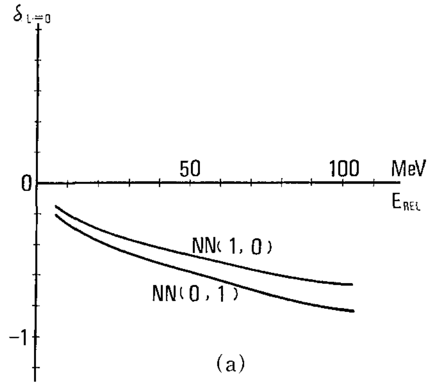
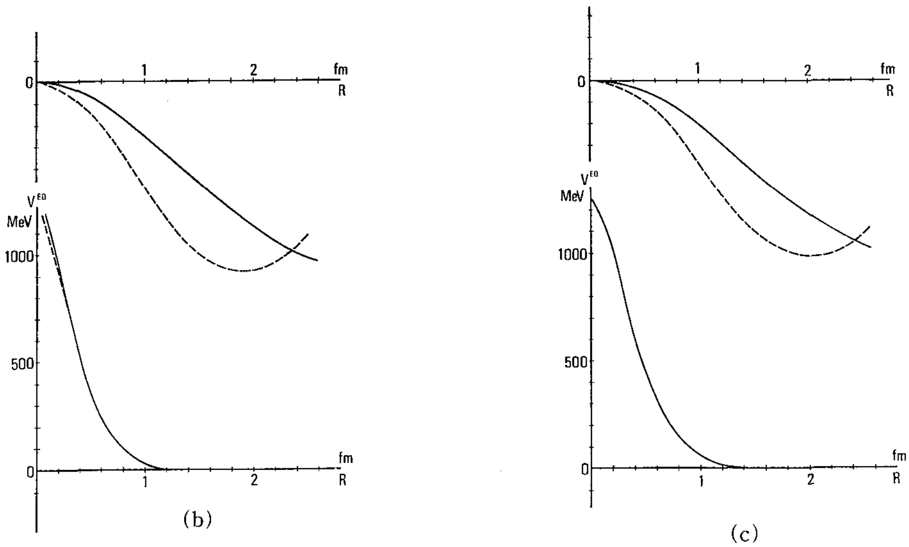


Fig. 5.(a) Phase shifts for the uncoupled (1, 0, *NN*) and (0, 1, *NN*) states.

(b) Wave functions and EQLP's for the uncoupled (1, 0, *NN*) state.

(c) Wave functions and EQLP's for the uncoupled (0, 1, *NN*) state.



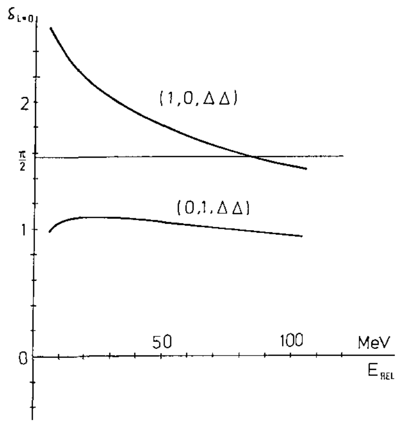


Fig. 6. Phase shifts for the uncoupled (1, 0, $\Delta\Delta$) and (0, 1, $\Delta\Delta$) states.

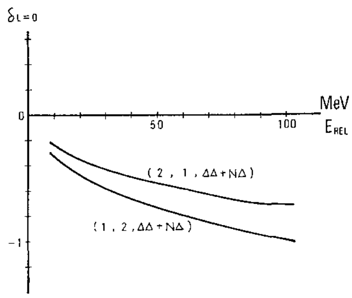


Fig. 7. Phase shifts for the coupled (2, 1, $\Delta\Delta$ $+N\Delta_a$) and (1, 2, $\Delta\Delta+N\Delta_a$) states.

tic) interaction is crucial in producing the short range repulsive force between two nucleons. The importance of the spin-spin interaction is also seen in the results for the uncoupled (1, 0, $\Delta\Delta$) and (0, 1, $\Delta\Delta$) states. Figure 6 shows that the $\Delta\Delta$ feel attractive forces in these states and we get the $\Delta\Delta$ bound state at $E_B = 30.5$ MeV for $(S, T) = (1, 0)$.

The calculations on the other symmetry mixed (2, 1, $\Delta\Delta + N\Delta_a$) and (1, 2, $\Delta\Delta + N\Delta_a$) states give the results similar to those for the above (1, 0) and (0, 1) cases. The phase shifts obtained by the coupled channel calculation are shown in Fig. 7. The repulsive cores at $R = 0.44$ and 0.48 fm for the (2·1) and (1, 2) states, respectively, are suggested from the behavior of the phase shifts.

3.5. *Dependence on the parameters*

Different choices of the parameters give no qualitative change to the previous results.

As stated in (I), the quadratic confinement potential does not contribute to the interaction between baryons and give the calculated phase shift which is *completely* independent of a'. Although this is not precisely the case for the linear confinement potential, it is also found to give little influence on the calculated phase shifts. Figure 8 compares the results for the coupled NN and $\Delta\Delta$ channels with two different values of a, i.e., $a = 141.2$ MeV/fm (set I) and $a = 0$. Such a drastic change of a gives rise to only a moderate change of the phase shifts. We therefore conclude that the mechanism of confinement does not significantly affect the interaction between two baryons within the framework of the present formulation.

The radius d of the repulsive core, defined by the average slope of the phase shift, changes significantly with the extension parameter b. For the ^3S and ^1S of the coupled $NN + \Delta\Delta$ channel, we obtain

$d = 0.67$ fm ($b = 0.8$ fm, param. set II in Table I),
0.39 fm ($b = 0.6$ fm, param. set I)

and

0.27 fm ($b = 0.5$ fm, param. set III)

for ^3S and

$d = 0.70$ fm (b $= 0.8$ fm, param. set II),
0.41 fm ($b = 0.6$ fm, param. set I)

and

0.31 fm ($b = 0.5$ fm, param. set III)

for ^1S.

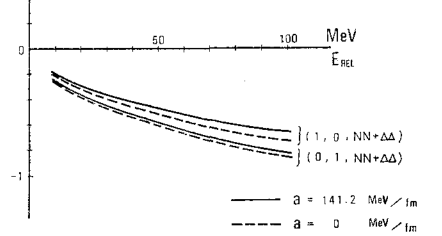


Fig. 8. Phase shifts for the coupled (1, 0, $NN + \Delta\Delta$) and (0, 1, $NN + \Delta\Delta$) states calculated for two different values of the strength of the linear confinement potential, that is $a = 141.2$ MeV/fm (full line) and 0 MeV/fm (dashed line).

The quark-gluon coupling constant α_s cannot be changed, especially, for the coupled channel calculations as α_s determines the N-Δ mass splitting and so the difference between the thresholds of the two coupled channels. However, the meson cloud may modify the mass splitting and therefore the actual value of α_s may be different from that determined by the mass splitting. In order to study the influence of this modification, calculations are done for the case where the strength of the "Coulomb" part of the potential is changed to twice or half of the initial value with the spin-spin term unchanged. Almost the same results are obtained. On the other hand, the modification of the strength of the spin-spin interaction gives a significant change on the calculated phase shifts, although it also changes the mass difference between N and Δ and therefore may not stand physical interpretations.

In conclusion, no qualitative changes are induced by the changes of the parameters, although the radii of the repulsive core are dependent on both the extension b of the quark wave function and the strength α_s of the spin-spin interaction.

3.6. *Effects of the effective meson exchange potential*

We have seen that the baryon-baryon interaction due to the quark exchange is mostly repulsive and cannot reproduce the attractive part of the nuclear force. Several other effects, however, are considered to contribute to the attractive interaction. They are

(1) the effect of the meson cloud surrounding the quark core, which is responsible for the meson exchange interaction,

(2) the core polarization effect or the change of the intrinsic wave functions ϕ_A and ϕ_B during the collision process, and

(3) the effect of tensor force in OGEP, which causes the mixing of D-state both in the intrinsic and scattering wave functions. The last one can be included within the framework of RGM but the effect on the scattering will not be too important since we are mainly considering the short range region where the mixing of the higher partial waves is expected to be small. The second one is certainly important but is also strongly dependent on the details of the dynamics, especially on the mechanism of confinement. It gives rise to, for example, the van der Waals force discussed previously, which should be considered as an unfavourable feature of the potential model. We have therefore decided not to include the core polarization effect and have restricted ourselves to the effects which are least dependent on the confinement mechanism.

Following the prescription given in § 5 of (I), we simulate the first one by introducing an effective meson exchange potential (EMEP) defined in Eq. (5·13) of (I) for the NN channels, where phenomenological potentials are known fairly well. The coupling between NN and $\Delta\Delta$ channels is neglected for simplicity, since the effect was found to be small except near the $\Delta\Delta$ threshold. A local potential of the Gaussian form is assumed for the EMEP, i. e.,

$$\mathcal{V}^{\text{EMEP}}(R) = V_0 \exp\left(-\frac{R^2}{a^2}\right). \tag{3·1}$$

We determine the strength V_0 and the range a by fitting the scattering length and the effective range for the ^3S and ^1S N-N states. The resulting values of V_0 and a are found to be the same for both states and are

$$V_0 = -394 \text{ MeV}$$

and

$$a = 0.94 \text{ fm}.$$

The calculated phase shifts, which are shown in Fig. 9, qualitatively reproduce the observed ones and indicate that the repulsive cores induced by the quark exchange effect are not smeared out by the attractive potential. The core radii, on the contrary, become larger than those without the EMEP, i. e.,

$$d = 0.67 \text{ fm} \quad \text{for} \quad {}^3S \quad \text{and}$$
$$0.60 \text{ fm} \quad \text{for} \quad {}^1S$$

for the parameter set I.

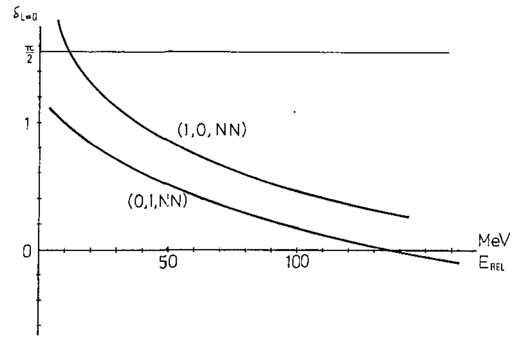


Fig. 9. Phase shifts for the uncoupled (1, 0, *NN*) and (0, 1, *NN*) states calculated with the effective meson exchange potential described in the text.

§ 4. Discussion and conclusion

The present calculation shows that the short range repulsion between two nucleons can be understood as a combined effect of the Pauli principle and the spin-spin interaction between quarks. The strong repulsion is not universal, however, and even attractive interactions are predicted in the $(S, T)=(3, 0)$, $(1, 0)$ and $(0, 1)\Delta\Delta$ states. With intermediate range attraction due to the meson exchanges, one expects deeply bound configurations in these states which may correspond to the "dibaryon resonances" suggested by Kamae et al.[9] In the case of $(S, T)=(3, 0)$, the relative wave function for such a bound state is expected to be similar to that of the $0s$ orbit in the harmonic oscillator model and, therefore, the bound state will be approximately described by the $(0s)^6$ configuration in the quark shell model. One may thus call it a 6 quark state rather than a two baryon state. For the $(S, T)=(1, 0)$ and $(0, 1)$ states, the coupling to NN channels is quite important and it is even not clear whether they remain as resonances when the coupling is introduced.

The symmetry properties of two baryon states have been considered by Neudatchin et al.[8] They introduced an attractive spin-isospin exchange force between quarks, which effectively prohibits the orbital [6] symmetry for the S wave NN states. They speculated that the NN relative wave function would have a node, the position of which corresponded to the radius of the "hard core", and also that the $(S, T)=(3, 0)$ and $(0, 3)$ states might have bound states. The spin-spin interaction in the present calculation plays essentially the same role as their spin-isospin exchange force, and their speculation is realized, although the symmetry between spin and isospin in their model disappears in our model.

Adiabatic approaches have been frequently employed to study nuclear force in quark models. Typical ones are Liberman's work[6] in a nonrelativistic quark model and DeTar's work[11] in the MIT bag model. They assume that the relative motion between two nucleons is slow compared with the internal motion of quarks in the nucleons and that the quantum fluctuation of the center of mass coordinates of the nucleons can be neglected. The latter seems to be a very strong assumption for a few-body system, and it is interesting to examine its validity. An advantage of the present approach is that one can treat the relative motion between two nucleons in a fully quantum mechanical way. Neglecting the fluctuation mentioned above corresponds to calculating the potential by the diagonal part of the GCM kernels in the present approach, which is defined by

$$V^{AD}(R)\equiv(K_{\mathrm{GCM}}^{(L=0)}(R, R)+V_{\mathrm{GCM}}^{(L=0)}(R, R))/N_{\mathrm{GCM}}^{(L=0)}(R, R) -(R\to\infty). \qquad (4\cdot1)$$

Figure 10 shows $V^{AD}(R)$ thus calculated together with the corresponding equivalent local potential $V^{EQ}(R)$ for the $(S, T)=(1, 0)$ and $(0, 1)$ states in the case where the $\Delta\Delta$ channels are omitted. The smearing effect due to the fluctuation is clearly seen and we can conclude that a quantal treatment of the center of mass motion is necessary for the quantitative discussion of the nucleon-nucleon interaction in any quark model.

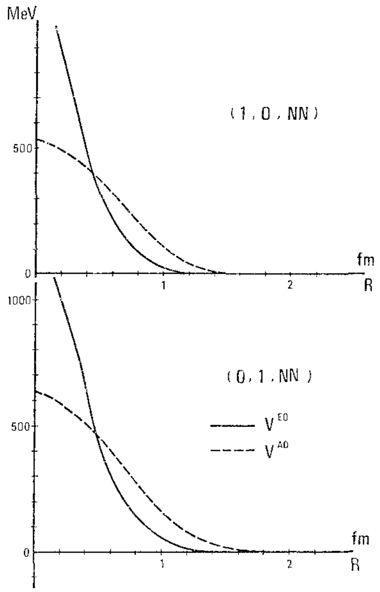


Fig. 10. The equivalent local potentials $V^{EQ}(R)$ (solid lines) and the adiabatic potentials $V^{AD}(R)$ (dashed lines) for the uncoupled (1, 0, NN) and (0, 1, NN) states. Parameters are those of the set I in Table III.

The adiabatic approach has also been employed by Harvey[12] in his recent paper on the quark model, where a part of the core polarization effect has been taken into account in the form of the "hidden color" channel. His conclusion is that the effect brings about strong attraction and washes out the short range repulsion for the NN channel. Although his model of quark dynamics is not the same as ours, qualitative features of his results will remain unchanged in the present model and the core polarization may appreciably alter the results of the present calculations. On the other hand, the effect is very sensitive to the confinement mechanism, and, as was discussed in (I), gives rise to the undesirable long range van der Waals force. The confinement described by the two body potential is therefore not appropriate for a dynamical treatment of the core polarization. In this respect, we note that our results are almost independent of the confinement term as mentioned in the last section and also in (I). Therefore we expect that the qualitative features of the present calculation will remain unchanged even if the confinement mechanism is replaced by an improved one which is presently unknown.

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