$(q\bar{q})$ AND $(q\bar{q})^2$ EXCITATIONS IN A QUARK MODEL OF THE NN INTERACTION

Y. FUJIWARA and K.T. HECHT

Physics Department, University of Michigan [1], Ann Arbor, MI 48109, USA

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The $(q\bar{q})$ and $(q\bar{q})^2$ excitations generated by off-shell terms of the Fermi–Breit quark–gluon interaction are incorporated into a quark model of the NN system leading to a semiquantitative fit of the NN scattering data.

In the past few years there have been many attempts to gain a more fundamental understanding of the NN interaction in terms of QCD-inspired quark models. The most detailed studies of the two-nucleon system have been carried out within the framework of the resonating group method through models in which a gluon exchange potential, usually in a one-gluon exchange approximation through the color analog of the Fermi–Breit interaction, is augmented by a phenomenological confining potential, (see ref. [1] for many earlier references). Such quark potential models may be open to question when used for the prediction of single baryon properties (e.g., absolute values of baryon masses) which depend on the nature and strength of the confining potential. Since the NN scattering phase shifts pass the crucial test of being almost completely insensitive to large changes in the strength of the confinement potential, the study of the NN interaction is free of such difficulties. In a recent study [1] the quark–antiquark excitations inherent in the quark–gluon interaction lagrangian have been explicitly incorporated into the quark model of the nucleon in order to study the effects of such excitations on the NN interaction. Quark exchange kernels for the two-nucleon system in which the three-quark (3q) components of the single-nucleon wave functions are augmented by (3q) $(q\bar{q})$ components lead to effective potentials with a medium-range attractive part and a greatly reduced repulsive core with the strong energy dependence and numerical values very similar to those of the short-range phenomenological terms of the Paris potential. By isolating those exchange terms which correspond to an exchange of a $(q\bar{q})$-pair between the two nucleons [2] it has also been possible to make good contact with conventional OBEP's through this extended quark model. The contribution of a particular (3q) $(q\bar{q})$ component of the single-nucleon wave function, for which the (3q) component has the quantum numbers of a nucleon and the $(q\bar{q})$ component the quantum numbers of a real pseudoscalar or vector meson, leads to an effective potential which is in remarkably good agreement with the corresponding OBEP for $R \gtrsim 1.2$ fm and has the same qualitative radial features over an even wider range. The simple $(q\bar{q})$ exchange potentials also have all the characteristics of convential OBEP's in their dependence on nucleon $(\boldsymbol{\sigma}_1 \cdot \boldsymbol{\sigma}_2)$ and $(\boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2)$ factors, and the relative importance and signs of spin–spin, spin-independent central, LS, and tensor terms. The major quantitative failure of the simple $(q\bar{q})$ exchange potential involves the pion tensor term which is too weak by a factor of ~3. This is in agreement with the predicted [1] pion–nucleon coupling constant, $g^2_{\mathrm{NN}\pi}$, which is also too weak by a factor of ~3. Since a simple $(q\bar{q})$ cluster with the quantum numbers of a pion cannot be expected to give a realistic picture of the pion, a quantitative fit of the OPEP (including its long-range Yukawa tail) was not expected. Nevertheless, with an adjustment of the pion tensor term to match the experimental strength, the extended quark model [1,2] gives a remarkably good account of spin–orbit

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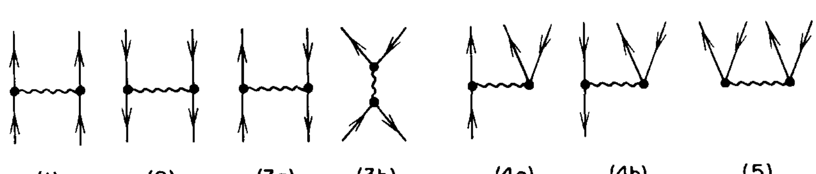
and tensor terms. Although the full central potentials gain a medium-range attractive part through the $(q\bar{q})$ excitations, this attractive part is too weak to bind the deuteron or fit the low energy phase shifts. Since the $(q\bar{q})$ excitations inherent in the quark–gluon interaction lagrangian cannot carry the quantum numbers of a scalar σ or δ meson, the counter-part of the conventional σ and δ meson exchange potentials was missing from the quark model of refs. [1,2]. However, excitations with the quantum numbers of a σ or δ meson can be incorporated into the extended quark model of the NN interaction through the $(q\bar{q})$ $(q\bar{q})$ excitations generated by RPA-type off-shell terms which are also a natural part of the full Fermi–Breit quark–gluon interaction. It is the purpose of this short note to show that the inclusion of such terms leads to quark exchange kernels which give a satisfying semiquantitative fit of the experimental NN scattering data.

Fig. 1 shows the five types of interaction terms of the Breit quark–gluon interaction. Off-shell contributions of the $(q\bar{q})$-pair creation type 4(a) were incorporated into the quark model of refs. [1,2]. Off-shell contributions of type (5) can lead to additional $(q\bar{q})$ $(q\bar{q})$-excitations in a multiquark system. As in refs. [1,2] improved single-nucleon wave functions are first calculated in which the dominant (3q) component is augmented by the (3q) $(q\bar{q})$ components generated by the interactions (4a) as well as the (3q) $(q\bar{q})$ $(q\bar{q})$ components generated by interactions of the type (5). Quark-exchange kernels are then calculated for the two-nucleon system described by these improved single-nucleon internal wave functions. The improved single-nucleon wave function now has the form

$$\phi_N = c_0\phi_0(3q) + \sum_{\alpha=1}^{24} c_\alpha\phi_\alpha((3q)(q\bar{q})) + \sum_{\beta=1}^{6} c_\beta\phi_\beta((3q)(q\bar{q})(q\bar{q})) , \qquad (1)$$

where the ϕ_α are defined through cluster RGM wave functions by eq. (12) of ref. [1]. In the ϕ_β it is useful to transform the $(q\bar{q})$ $(q\bar{q})$ pieces into $(q)^2$ $(\bar{q})^2$ form [3], where the (3q) component with the color, spin, isospin of a real nucleon is coupled to $(q)^2$ with S_1T_1 and matching color symmetry which is first coupled to $(\bar{q})^2$ with S_2T_2 where the $(q)^2$ $(\bar{q})^2$ cluster is assumed to be coupled to the color singlet, S, T quantum numbers of σ or δ type. All internal and relative motion functions in ϕ_β are assumed to be 0s oscillator functions. The kernels coupling ϕ_β to ϕ_0 and ϕ_β to $\phi_{\beta'}$ have been evaluated by the methods of ref. [1]. Only exchange terms are retained for ϕ_0 to ϕ_β coupling with $S = 0$, $T = 0$, under the assumption that the disconnected diagrams of the direct term should be excluded for a real nucleon as for the real vacuum. Diagonalization in the full 31-dimensional space of eq. (1) then leads to a new set c_0, c_α, c_β. The β with $S = 0$, $T = 0$ and S_1T_1, S_2T_2 of 00,00; 11,11; 01,01; and 10,10 are denoted by σ_1; σ_2; σ_3; and σ_4. For $S = 0$, $T = 1$ the states with S_1T_1, S_2T_2 of 11,11 and 01,01 are denoted by δ_1 and δ_2. The resultant c_β are: $c_{\sigma_1} = -0.004$, $c_{\sigma_2} = -0.100$, $c_{\sigma_3} = -0.195$, $c_{\sigma_4} = 0.043$, $c_{\delta_1} = -0.053$, and $c_{\delta_2} = -0.138$.

The new c_α and c_0 are almost identical to those quoted in table 1 of ref. [1]. (E.g., $c_{N\omega 3/2}$ is changed to -0.216 from -0.232, whereas c_0 is changed to 0.849 from 0.857). Due to these small changes, and also due to the fact that the ϕ_β do not make a direct



$$H_{int} = H_{qq} + H_{\bar{q}\bar{q}} + (H_{q\bar{q}} + H_{q\bar{q}}^{(ann)}) + \left(H_{q\to qq\bar{q}} + \text{c.c.}\right) + \left(H_{\bar{q}\to\bar{q}q\bar{q}} + \text{c.c.}\right) + \left(H_{\to(q\bar{q})^2} + \text{c.c.}\right)$$

Fig. 1. The full Breit interaction hamiltonian.

contribution to the single-nucleon electromagnetic properties and the nucleon vector-meson coupling constants, no attempt was made to reevaluate the four parameters of the model. The values of ref. [1] were retained for α_s (strong coupling constant), m (quark mass), b (oscillator length parameter), and a_c (confinement potential constant). Although the improved single nucleon wave function no longer gives a fit to the N and Δ masses, it was shown [1] that the effective NN potentials were given by the strengths of the c_α's and were not sensitive to the details of the parameter fit. It should also be noted that the ϕ_β, like the ϕ_α, form a nonorthogonal set. The largest overlap, (between the normalized ϕ_{σ_3} and $\phi\delta_2$) of −0.408 is now somewhat larger due to the pure 0s character of the ϕ_β. For the same reason the fully antisymmetrized ϕ_β of scalar meson (σ, δ) type also effectively carry some axial vector meson contributions.

The improved single nucleon functions of eq. (1) are used to evaluate the exchange kernels for the two-nucleon system

$$G(\mathbf{R}, \mathbf{R}') = c_0^4 G_0(\mathbf{R}, \mathbf{R}') + c_0^3 \left(\sum_{\alpha=1}^{24} c_\alpha G_\alpha(\mathbf{R}, \mathbf{R}') + \sum_{\beta=1}^{6} c_\beta G_\beta(\mathbf{R}, \mathbf{R}') \right) + \dots . \tag{2}$$

As in ref. [1] contributions of second order in the c_α and c_β are neglected. It should be noted that the kernels G_0 and G_α give rise to spin–orbit and tensor as well as central terms in the NN interaction, but the coupling kernels G_β, connecting the (3q)–(3q) $(q\bar{q})^2$ components to the (3q)–(3q) components of the NN system, make contributions only to the central terms. The kernels of eq. (2) are converted into equivalent local potentials through the Wigner transform-WKB approximation as in ref. [1]. In addition, the RGM equations have been solved directly. Fig. 2 shows the S-wave equivalent local potentials. The overall attraction for the $^1S_0$ potential at the low energy limit, $E_{cm} = 0$, comes from the combined effects of the (3q) $(q\bar{q})$ and (3q) $(q\bar{q})^2$ coupling kernels G_α and G_β which contribute −281 and −448 MeV, respectively at $R = 0$, overcoming a repulsive core of +669 MeV from the pure (3q)–(3q) kernel, G_0. The strongest energy dependence of these potentials arises through the (3q) $(q\bar{q})$ component. The weaker attraction for the low energy $^3S_1$ central potential is in accord with the experimental data since the binding in the

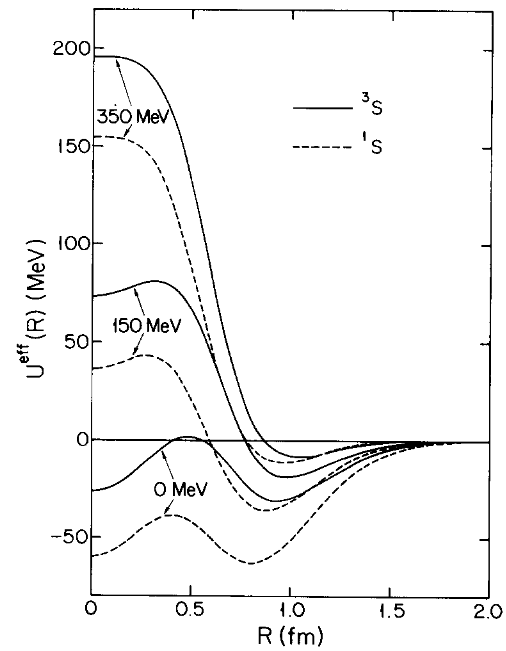


Fig. 2. The S-wave equivalent local potentials.

$^3S_1$ state gets important contributions through coupling to the $^3D_1$ channel via tensor force terms.

Fig. 3 shows the predicted RGM phase shifts for

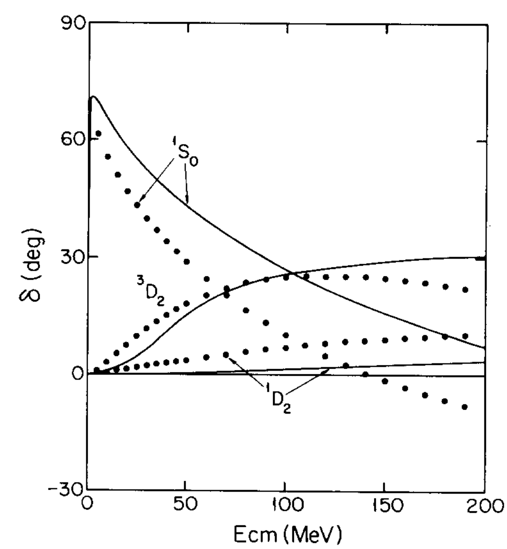


Fig. 3. Phase shifts given by single channel RGM calculations. The (dotted) experimental points for figs. 3–6 are taken from ref. [4].

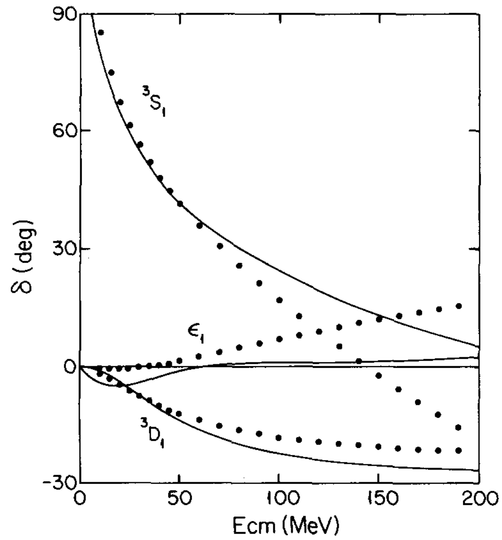


Fig. 4. The $^3S_1$ and $^3D_1$ phase shifts and mixing parameter. The predicted $^3S_1$ binding energy is 0.5 MeV.

the partial waves given by single-channel calculations. Although the predicted $^1S_0$ phase shifts are somewhat too attractive, particularly at higher energies, the $^1S_0$ state is unbound. It should also be borne in mind that this is a zero parameter prediction since the four parameters of our model, determined from the single-baryon data in ref. [1], were used without further adjustment. Fig. 4 shows the predicted $^3S_1$ and $^3D_1$ phase shifts including the important effects of channel coupling through the tensor force. (The $^3S_1$ phase shift predicted by purely central terms at E_{cm} = 20 MeV, e.g., would have been $\delta(^3S_1) = 19°$ as compared with the full prediction of 64° shown in fig. 4). Since the tensor component of the ($q\bar{q}$)-exchange potential of pion type predicted by our quark model is too weak by a factor of ~3 compared with conventional OPEP's, (as shown in fig. 10 of ref. [2]), this particular term was replaced with a more realistic OPEP, with coupling constant fitted to the experimental value, ($g_\pi^2(0) = 14.17$), and a momentum dependence of $g_\pi(k)$, as given by eq. (69) of ref. [1], approximated by a simple gaussian form. It is an advantage of our quark model that such specific terms in the predicted NN interaction can be isolated and improved. With the exception of this adjustment, however, the even partial-wave predictions of figs. 3 and 4 are free of parameter fitting, including an unadjusted

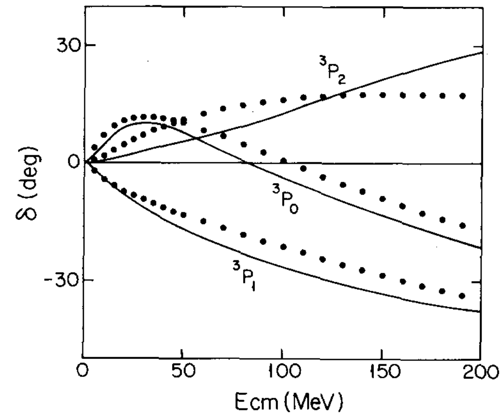


Fig. 5. The ^3P phase shifts with adjusted E_{internal} (see text).

predicted internal energy of 2(693 MeV).

Although comparable predictions for odd partial waves are not unreasonable, it is clear that the central terms of our odd-L potentials are somewhat too repulsive largely through contributions of the coupling kernels G_α of (3q) ($q\bar{q}$) type. To compensate for this repulsion the internal energy in the RGM equations for the odd partial waves has been adjusted to a new value of 2(250 MeV). The predicted internal energy is sensitive to the details of our model parameters. The central terms of the odd–L potentials may gain significant contributions from the neglected second

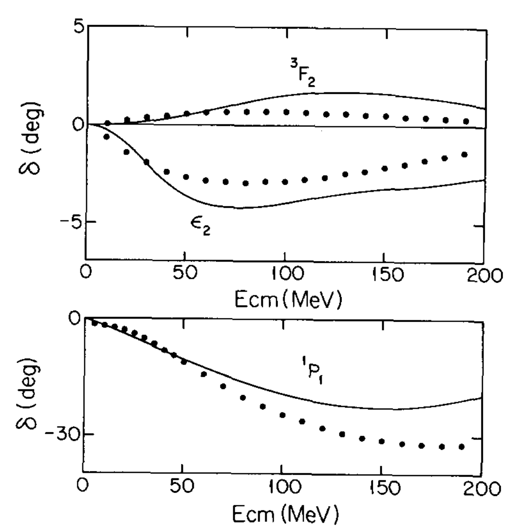


Fig. 6. The $^3F_2$ and $^1P_1$ phase shifts.

order terms of eq. (2) and can perhaps also not be expected to fit the needed experimental values. Since E_{internal} affects the RGM results through the norm kernel term, it can be adjusted to compensate partly for the central odd-L repulsion. With this one adjusted parameter, however, the odd-L phase shifts are fitted well by our model, as shown in figs. 5 and 6. The ^{3}P phase shifts, in particular, indicate that the predicted spin–orbit potentials have the right magnitude. It is interesting to note that about 65% of the triplet odd LS potentials in the $R \sim 1$ fm range arise from the coupling kernels, G_α, of (3q) (q$\bar{\mathrm{q}}$) type and only 35% from the (3q)–(3q) kernel, G_0, including both symmetric and antisymmetric LS terms [5].

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