Single pole dominance in short- and intermediate-range $NN$ interaction

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It is demonstrated that both elastic and inelastic $NN$ scattering at laboratory energies up to 600–800 MeV, at least in some partial waves characterized by a large inelasticity, can be described by a superposition of the conventional long-range one-pion exchange and a specific short-range interaction induced by the $s$-channel dibaryon exchange. For the $^3P_2$, $^1D_2$ and $^3F_3$ partial waves, the pole parameters giving the best fit of the real and imaginary parts of the $NN$ phase shifts are consistent with the parameters of the respective isovector dibaryon resonances found experimentally. In the $^1S_0$ channel, the suggested interaction gives two poles of the $S$-matrix — the well-known singlet deuteron and an excited dibaryon. On the basis of the results presented, a conclusion is made about the nature of $NN$ interaction and its strong channel dependence.

I. INTRODUCTION

Nowadays the Effective Field Theory (EFT), or Chiral Perturbation Theory (ChPT), is a dominating framework for the quantitative description of $NN$ interaction at low and moderate energies [1–4]. In this approach, a peripheral part of $NN$ interaction is described via a superposition of terms of the perturbation theory series corresponding to subsequent orders, i.e., the leading order (LO), the next-to-leading order (NLO), the next-to-next-to-leading order (N$^2$LO), etc., while the short-range contributions are parameterized through the so-called contact terms which, according to the general concept, should be independent of energy and the expansion order. By construction, this general approach is valid until the collision energies $E_{	ext{lab}} \approx 350$ MeV, while at higher energies it should be supplemented by an appropriate theoretical model to describe the short-range components of $NN$ interaction and the short-range $NN$ correlations in nuclei. For this purpose, one can consider the well-known quark model in its various versions (see, e.g., [2]). However such a hybrid approach inevitably leads to serious difficulties with double counting because in quark models the gluon exchanges between quarks are usually supplemented by meson ($\pi$ and $\sigma$) exchanges, which immediately results in appearance of not only short-range but also long-range meson-exchange forces between nucleons.

At the moment in nuclear physics there is a wide class of phenomena where one observes close connection between short-range correlations of nucleons in nuclei and some distortions of quark momentum distributions in the deuteron, $^3$He and other nuclei [6, 8]. These phenomena include the EMC effect, DIS observations, cumulative effects, etc. It is evident that neither the traditional meson-exchange nor the modern EFT approaches are relevant for description of such effects occurring at very high momentum transfers. On the other hand, since all these phenomena are closely related to nuclear force and nuclear structure at short distances, they should be described within some general scheme which includes the correct treatment of $2N$ and $3N$ forces.

Therefore, for further progress in this area it would be highly desirable to treat the short- and intermediate-range $NN$ interactions by using some QCD-motivated model which could reproduce correctly the basic effects of the nucleon quark structure in different $NN$ partial waves, however without addressing the all complexity of multi-quark dynamics. Such a model description of the short-range $NN$ dynamics should not modify the known interaction at long distances.

In our opinion, suitable objects, which, on the one hand, can reproduce the main features of six-quark dynamics and, on the other hand, are strongly coupled to the hadronic ($NN$, $N\Delta$ and $\Delta\Delta$) channels at low and moderate energies are dibaryon resonances which were predicted by Dyson and Xuong in 1964, at the very beginning of quark era [8]. Just recently, after many years of rejection, doubt, and contradictory findings, a number of dibaryon resonances have been eventually confirmed in the modern high-precision experiments [9–12] (see also the recent review [13]).

Dibaryon resonances are very attractive to describe the short-range $NN$, $N\Delta$ and $\Delta\Delta$ forces not only due to their six-quark structure but, first of all, because they are specific relatively long-lived states in which six-quark dynamics should be manifested most clearly. So, the goal of the present paper is to demonstrate that $NN$ interaction in many partial-wave channels can be described properly by a superposition of the long-range meson-exchange potentials and one simple $s$-channel exchange potential driven by a dibaryon pole.

For particular $NN$ partial-wave channels considered below (and for many other channels) it is crucial to include the $NN \rightarrow N\Delta$ (and/or $NN \rightarrow \Delta\Delta$) coupling. In turn, the $NN \rightarrow N\Delta$ transition may or may not be accompanied by a dibaryon formation. While the latter type of coupling has been effectively included into our model consideration through the resonance parameters
(in fact, the isovector dibaryon width is mainly due to the $D \to N\Delta$ decay), the former (pure $t$-channel) coupling is actually a background to the $s$-channel dibaryon generation which should be included explicitly into the $NN$ interaction potential. Though the background processes may affect significantly the description of $NN$ inelasticities above the resonance energy, they will not change our main results concerning the impact of the basic dibaryon mechanism at lower energies. Thus we postpone the consistent treatment of the $NN \to N\Delta$ coupling to the future work.

It should be noted that the initial version of the dibaryon model for $NN$ and $3N$ interactions developed in [14] (see also [15] and references therein) provided quite encouraging results in description of $NN$ scattering in various partial waves. In particular, it allowed to reproduce real parts of phase shifts in different $NN$ channels up to energies 500–600 MeV (lab.). However, since then dibaryons have achieved a more reliable experimental status, while the parameters of the real dibaryon poles found in the work [14] were never compared to those of the physical dibaryon resonances. So, in the present paper we tried to incorporate into the initial model the experimentally found dibaryon resonances together with their empirical parameters to describe simultaneously the real and imaginary parts of the $NN$ phase shifts in a broad energy range. The results of this study would allow to judge about the true applicability and merit of such a non-conventional description of $NN$ interaction.

The structure of the paper is as follows. In Sec. II we introduce the two-channel formalism with one external and one inner channel which is used further to describe the $NN$ partial phase shifts. In Sec. III and IV we present the results for the particular isovector $NN$ partial waves and illustrate effectiveness of the dibaryon mechanism. We summarize the results and conclude in Sec. V.

II. $NN$ SCATTERING DRIVEN BY A SINGLE STATE WITH A COMPLEX ENERGY IN THE INTERNAL CHANNEL

We consider below a two-channel model with one complex pole in the effective interaction potential. This model corresponds to the physical pattern of $NN$ scattering driven by the traditional one-pion exchange in the external channel and by one state with a complex eigenvalue (i.e., the “bare” dibaryon state) in the internal channel. The complex energy of this state can be interpreted as a consequence of different modes of its decay which do not include the $NN$ mode.\footnote{The initial dibaryon state with the real energy can be treated in terms of the field theory as a “bare” dibaryon, while the coupled-channel dibaryon which is able to decay into the $NN$ continuous spectrum can be identified as a “dressed” dibaryon. In this sense, the dibaryon with the complex energy can be called as a “semidressed” one. However, to retain the unified notation, we will refer to the initial dibaryon with the complex energy as a “bare” one.}

The total matrix Hamiltonian for such a two-channel problem has the form:

$$H = \left( \begin{array}{cc} h_{NN} & \lambda_1 |\phi\rangle \langle \alpha| \\ \lambda_1 |\phi\rangle \langle \alpha| & E_D |\alpha\rangle \langle \alpha| \end{array} \right),$$

where the external-channel Hamiltonian $h_{NN}$ acts in the space of $NN$ variables and includes the peripheral $NN$ interaction which is exhausted by the one-pion exchange potential (OPEP).\footnote{For definiteness, we choose the dipole form factor in the $\pi NN$ vertex with the soft cutoff parameter $\Lambda_{\pi NN} = 0.65$ GeV.} In case of the single-pole model, the Hilbert space of the inner channel is one-dimensional and therefore the internal Hamiltonian is reduced to a single term with a complex eigenvalue $E_D = E_0 - i \Gamma_D / 2$.

To determine the form factor $|\phi\rangle$ of the transition between the external ($NN$) and the internal (dibaryon) channels, it is necessary to use some microscopic model that describes both channels within a unified approach. In our previous works [14, 15], a dibaryon model for nuclear forces was developed, in which a microscopic six-quark shell model in a combination with the $3P_0$ quark mechanism of pion production was used to determine the transition amplitude between the channels. Note that the transition form factor $|\phi\rangle$ is a function of the relative coordinate $r$ or the relative momentum in the $NN$ channel, and also depends on the spin, isospin, orbital and total angular momenta of the two-nucleon state. Here we use the same form of this function as in Ref. [14], i.e., the harmonic oscillator term:

$$\phi(r) = N r^{l+1} \exp \left[ -\frac{1}{2} \left( \frac{r}{r_0} \right)^2 \right],$$

where $l$ is the orbital angular momentum in the two-nucleon system, $N$ is a normalization factor, and $r_0$ is a scale parameter.

Since the total Hamiltonian $H_{\text{eff}}$ couples the inner channel with the $NN$ relative-motion channel, it is convenient to exclude the internal channel by the commonly-used method [10] and to deal with the $NN$ relative motion only. Then the respective effective energy-dependent Hamiltonian in the $NN$ channel takes the form:

$$H_{\text{eff}}(E) = h_{NN} + \frac{\lambda_1^2 |\phi\rangle \langle \phi|}{E - E_D}.$$  

Due to the fact that the basic term of the effective Hamiltonian $h_{NN}$ has a separable form, it is convenient to define explicitly an additional $t$-matrix in the distorted-wave representation:

$$t(E) = \frac{\lambda_1^2 |\phi\rangle \langle \phi|}{E + i0 - E_D - J_1(E + i0)},$$

The initial dibaryon state with the real energy can be treated in terms of the field theory as a “bare” dibaryon, while the coupled-channel dibaryon which is able to decay into the $NN$ continuous spectrum can be identified as a “dressed” dibaryon. In this sense,
where $J_1(Z)$ is the matrix element of the resolvent of the external $NN$ Hamiltonian $g_{NN}(Z) \equiv [Z - h_{NN}]^{-1}$:

$$J_1(Z) = \lambda_0^2 \langle \phi | g_{NN}(Z) | \phi \rangle.$$

(5)

Note that the imaginary part of this function at a real positive energy can be found in an explicit form as

$$\text{Im} \ J_1(E + i0) = -\pi \lambda_0^2 |\langle \phi | \psi_0(E) \rangle|^2,$$

(6)

where $|\psi_0(E)\rangle$ is the scattering function for the external Hamiltonian $h_{NN}$. Using the standard formula for the transition operator, one can easily obtain the expression for the total $S$-matrix:

$$S(E) = e^{2i\delta_0} \frac{E - E_D - J_1(E + i0)}{E - E_D - J_1(E + i0)},$$

(7)

where $\delta_0(E)$ is the phase shift for the external Hamiltonian $h_{NN}$. Thus, the pole of the total $S$-matrix can be found in a complex energy $Z$ plane from the condition:

$$Z - E_D - J_1(Z) = 0.$$

(8)

The condition makes it possible to find the renormalized position of the dressed dibaryon resonance (relatively to the initial “bare” value of $E_D$), i.e., the complex function $J_1(Z)$ causes shifts for the real and imaginary parts of the dibaryon pole. These shifts arise due to coupling of the initial dibaryon to the external $NN$ channel.

Thus, we have formulated a simple model for coupling between the external (OPEP) $NN$ channel and the internal (“bare” dibaryon) channel which leads to a renormalization of the complex energy of the initial “bare” dibaryon and its transformation to the real mass and width of the “dressed” dibaryon. The simple mechanism for coupling between the external and internal channels can be represented graphically by the diagram series shown in Fig. 1. This series corresponds to the well-known Dyson equation for a dressed particle in the quantum field theory.

![Diagram](image)

FIG. 1: The sum of terms corresponding to dressing of the total dibaryon propagator. $D_0$ and $D$ are the propagators of the “bare” and “dressed” dibaryons, respectively.

### III. $NN$ Scattering in Isovector Channels

In this section, we consider as particular examples the isovector partial waves ($^3P_2$, $^1D_2$ and $^3F_3$) of $NN$ scattering, where the empirical phase shifts do not manifest explicitly the behaviour inherent to the repulsive core, at least until relatively high collision energies $E_{lab} \approx 800 \text{ MeV}$. The resonance behaviour in these partial waves near the $N\Delta$ threshold was established long ago in experiments $^{17,18}$ and was interpreted either by the threshold effect or the true dibaryon formation. The dibaryon interpretation was then supported by the partial-wave analyses of different groups which found the $S$-matrix poles corresponding to dibaryon resonances $^{21,22}$ (see also the review paper $^{23}$). In our recent works $^{24,25}$ the importance of these three resonances was established for description of the basic pion-production reaction $pp \rightarrow d\pi^+$, and the $^3P_2$ dibaryon was shown to play a crucial role in reproducing the polarization observables. Simultaneously, the $^3P_2$ dibaryon was observed in the recent high-precision experiment on the reaction $pp \rightarrow pp(^3S_0)\pi^0$ $^{11}$.

It should be emphasized that previous analyses of $NN$ scattering which considered dibaryon resonances generally suggested that the dibaryon pole can give a significant contribution to the $NN$ phase shifts in a respective partial wave only near the resonance energy. Contrary to this, we will show below that the single pole (combined with a long-range OPEP contribution) can explain the phase shifts behaviour in the isovector channels in a broad energy range from zero up to the resonance energy. The isoscalar channels of $NN$ scattering will be treated in our subsequent paper.

#### A. Channel $^3P_2$

The empirical $NN$ phase shifts in the triplet channel $^3P_2$ as found by the George Washington University group (SAID) $^{26}$ do not display any sign of the repulsive core, at least up to energy of 1000 MeV (lab.), and remain to be positive in this energy range. $^4$ This can be interpreted as a fact that the traditional repulsive core does not play a crucial role in this channel, and thus $NN$ interaction here is managed by a rather strong attraction, which is likely due to generation of a dibaryon resonance.

Below we show that such an attraction can be reproduced by a single dibaryon pole in the effective Hamiltonian $h_{NN}$ via varying its position $E_D = E_0 - i\Gamma_D/2$ and the coupling constant $\lambda_1$ for the external and internal channels. To take into account the unitarity condition at energies below opening of inelastic channels in $NN$ interaction, we included the energy dependence of the bare

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$^3$ The fact that at higher energies $E_{lab} > 800$ MeV the real phase shifts in some channels become negative can be explained by strong absorption at these energies, i.e., by appearance of a large imaginary part of phase shifts that can be interpreted as a consequence of strong repulsion.

$^4$ Such a behaviour can be explained in the conventional $NN$ potential models by a complete compensation of the short-range repulsive core by the very strong attractive spin-orbital potential, so that the resulting potential in this channel turns out to be negative $^{27}$.  

dibaryon width $\Gamma_D$ by using a smoothed step-like shape:

$$\Gamma_D(E) = \Gamma_0 \left[1 + \exp \left(\frac{E_{\text{thr}} - E}{\alpha}\right)\right]^{-1}. \quad (9)$$

Here $\alpha$ is a smoothing scale parameter and the energy $E_{\text{thr}} = m_\pi + \beta$ is related to the pion-production threshold, where $m_\pi$ is the averaged pion mass and $\beta$ is some energy shift.

The comparison of empirical (SAID) and theoretical $^3P_2$ $NN$ phase shifts in the energy interval $E_{\text{lab}} = 0$–$800$ MeV is presented in Fig. 2 (real part) and Fig. 2b (imaginary part).\(^5\) The parameters for the potential have been taken as $\lambda_1 = 0.06$ GeV and $r_0 = 0.71$ fm. The initial dibaryon position corresponds to $M_D = 2.195$ GeV and $\Gamma_0 = 0.074$ GeV (here $M_D = E_0 + 2m_N$, where $m_N$ is the nucleon mass).

FIG. 2: Real (a) and imaginary (b) parts of the partial $NN$ phase shifts in the $^3P_2$ channel (solid curves) in comparison with the SAID data [26] (points).

It is seen from Fig. 2 that the single-pole model in a combination with a simple OPEP provides almost quantitative agreement with the empirical data for the real part and reasonable agreement for the imaginary part of the $^3P_2$ phase shifts up to energies $E_{\text{lab}} \simeq 700$ MeV. Needless to say that the above agreement for both real and imaginary parts of the phase shifts, i.e., for elastic $NN$ scattering and meson production simultaneously, was attained with the same parameters for the bare dibaryon. Some discrepancy observed for the imaginary part of the phase shift, starting just above the pion-production threshold, is likely related to our simplified description of the energy dependence of the dibaryon width, and can probably be decreased by a more sophisticated treatment of $\Gamma_D(E)$.

The proper way to proceed here is to include explicitly in the model the $N\Delta$ $P$-wave channels which couple to the $NN(^3P_2)$ channel.

Now, using Eq. (9), one can find easily the renormalized position of the resonance pole in this channel. We found the following parameters for the dressed dibaryon in the channel $^3P_2$: $M_{\text{th}}(^3P_2) = 2.22$ GeV and $\Gamma_{\text{th}}(^3P_2) = 0.095$ GeV. These parameters should be compared with the respective experimental values found by the ANKE-COSY Collaboration [11]: $M_{\text{exp}}(^3P_2) = 2.197(8)$ GeV and $\Gamma_{\text{exp}}(^3P_2) = 0.130(21)$ GeV (the numbers in parentheses denote the error in the last digit).\(^7\) Thus, the mass and width obtained in this work for the $^3P_2$ dibaryon turn out to be rather close to their experimental values (with account of the quoted errors).

Since our model does not include any other assumptions, except the OPEP and presence of a single dibaryon pole, we owe to conclude that the results obtained point directly to the dominance of the intermediate dibaryon production in the $NN$ interaction in this channel.

B. Channel $^1D_2$

A fully similar consideration for $NN$ phase shifts in the singlet channel $^1D_2$ leads to theoretical predictions shown in comparison with the respective empirical data (SAID) in Figs. 3a and 3b for the real and imaginary parts of the phase shifts, respectively. The potential parameters are $\lambda_1 = 0.045$ GeV, $r_0 = 0.79$ fm. The initial dibaryon position is $M_D = 2.16$ GeV and $\Gamma_0 = 0.074$ GeV.

FIG. 3: Real (a) and imaginary (b) parts of the partial $NN$ phase shifts in the $^1D_2$ channel (solid curves) in comparison with the SAID data [26] (points).

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5 For all cases discussed in this section, the values $\alpha = 0.02$ GeV and $\beta = 0.02$ GeV have been used.

6 Though the SAID results for $np$ scattering are plotted in Fig. 2 and subsequent figures, the $pp$ scattering phase shifts are indistinguishable from them in the scale of the picture.

7 The authors of the work [11] used two different fitting procedures, the first leading to the above quoted values for the mass and width of the $^3P_2$ dibaryon and the second — to somewhat larger values still consistent with the above ones within errors.
Here we see again almost perfect agreement for the real part and reasonable agreement for the imaginary part of the phase shifts at energies up to 600 MeV. The parameters of the dressed dibaryon in the $^{1}D_{2}$ channel found in our calculations are $M_{\text{th}}(^{1}D_{2}) = 2.175$ GeV and $\Gamma_{\text{th}}(^{1}D_{2}) = 0.104$ GeV. These parameters are consistent with those found previously in experiments [17, 19] and partial-wave analyses [20, 22] and turn out to be particularly close to the values obtained in the early experiment of 1955 [17] where the $^{1}D_{2}$ resonance was first observed: $M_{\exp}(^{1}D_{2}) \simeq 2.16$ GeV and $\Gamma_{\exp}(^{1}D_{2}) \simeq 0.12$ GeV.

It is important to emphasize that the $NN$ channel $^{1}D_{2}$ strongly couples to the $N\Delta$ channel $^{3}S_{2}$, and just this strong coupling with a subsequent decay of the $\Delta$-isobar determines a large portion of inelasticity in the $^{1}D_{2}$ partial wave. Since we did not take into account the t-channel $\Delta$ excitation in our model, the description of inelastic phase shifts turns out to be not perfect. However, almost quantitative description of the real part of the phase shifts until $E_{\text{lab}} = 600$ MeV and reasonable description of their imaginary part (with the same model parameters) point to the dominant role of the dibaryon resonance (including its $N\Delta$ component) in this channel as well.

C. Channel $^{3}F_{3}$

First experimental evidence of the $^{3}F_{3}$ dibaryon resonance with the mass $M_{D} \simeq 2.26$ GeV was obtained in 1977 [18]. It was also established that there is a very large inelasticity in the $^{3}F_{3}$ $NN$ partial wave. So, an interesting and challenging question arises: to what degree the $^{3}F_{3}$ resonance contributes to the real and imaginary parts of the $NN$ phase shifts in this partial channel. From the first glance, the possible impact of the dibaryon can be seen mainly at the energies close to the resonance position. However, in general, the influence of the s-channel dibaryon exchange in $NN$ scattering must be retraced also far from the resonance energy and give an evident impact to the off-shell t-matrix.

In Fig. 4 the real and imaginary parts of the partial phase shifts in the $^{3}F_{3}$ channel are compared with the empirical (SAID) data. As is clearly seen from the figure, the single $^{3}F_{3}$ pole (in a combination with OPEP) can reproduce very well the real $NN$ phase shifts from zero energy up to 700 MeV (see Fig. 4).

The imaginary part of the phase shifts turns out to be underestimated in this case. This underestimation is likely related to coupling between the $N\Delta(^{2}P_{3})$ and $NN(^{3}F_{3})$ states which was not consistently treated in our present model. It should be also mentioned that the phase shift $\delta_{0}$ corresponding to the initial OPEP in the external $NN$ channel though being very small depends strongly on the value of the cutoff parameter $\Lambda$. So that, different fits can be obtained for different $\Lambda$ values as is shown in Fig. 4. However, for consistency, one should choose the resonance parameters corresponding to the same value $\Lambda = 0.65$ GeV as was used for description of the lower partial waves. The parameters of the coupling potential in this case are $r_{0} = 0.55$ fm$^{-1}$ and $\lambda_{1} = 0.06$ GeV, while the pole parameters are $M_{D} = 2.24$ GeV and $\Gamma_{0} = 0.12$ GeV.

The parameters of the dressed dibaryon found for the $^{3}F_{3}$ channel, i.e., $M_{\text{th}}(^{3}F_{3}) = 2.226$ GeV and $\Gamma_{\text{th}}(^{3}F_{3}) = 0.156$ GeV, are also consistent with previous findings [18–22] and occur to be rather close to the values obtained in the first experimental observation of the $^{3}F_{3}$ resonance [18]: $M_{\exp}(^{3}F_{3}) \simeq 2.26$ GeV and $\Gamma_{\exp}(^{3}F_{3}) \simeq 0.2$ GeV.

The results presented in this section for the isovector dibaryon resonances together with experimental data are summarized in Table I.

![FIG. 4: Real (a) and imaginary (b) parts of the partial $NN$ phase shifts in the $^{3}F_{3}$ channel for different values of the OPEP cutoff parameter $\Lambda = 0.85$ GeV (dash-dotted curves) and 0.65 GeV (solid curves) in comparison with the SAID data [26] (points).](image)

| $^{2S+1}L_{J}$ | $M_{D}$ | $\Gamma_{0}$ | $M_{\text{th}}$ | $\Gamma_{\text{th}}$ | $M_{\exp}$ | $\Gamma_{\exp}$ |
|---------------|--------|-------------|-----------------|-------------------|------------|----------------|
| $^{3}P_{2}$   | 2.195  | 0.074       | 2.220           | 0.095             | 2.197(8)   | 0.130(21)     |
| $^{1}D_{2}$   | 2.160  | 0.074       | 2.175           | 0.104             | 2.14–2.18  | 0.05–0.1      |
| $^{3}F_{3}$   | 2.240  | 0.120       | 2.226           | 0.156             | 2.20–2.26  | 0.1–0.2       |

By using the developed model, we can also readily estimate the elastic ($NN$) part of the total dibaryon width by simply comparing the values of $\Gamma_{0}$ and $\Gamma_{\text{th}}$. Thus, the decay width into the $NN$ channel turns out to be about 22%, 29% and 23% of the total width for the $^{3}P_{2}$, $^{1}D_{2}$ and $^{3}F_{3}$ resonances, respectively. These values agree qualitatively with the previous estimates [21, 22] which typically give the $D \to NN$ partial width to be 10–20% of the total width for the isovector dibaryons.
Thus, in this section we have got a very good quantitative approximation for the real parts of the partial phase shifts in the \( N\bar{N} \) channels considered and a rather good qualitative approximation for the imaginary parts of these phase shifts in a broad energy range starting from zero energy which is very far from the position of the “bare” dibaryon. The crucial point of the results presented above is an agreement of the masses and widths of the dressed dibaryons obtained by fitting the \( N\bar{N} \) phase shifts in our model with the parameters of experimentally found dibaryons.

IV. CHANNEL \(^1S_0\), DESCRIPTION OF THE REPULSIVE CORE EFFECTS

One of the main ingredients of the conventional models for \( N\bar{N} \) interaction is a well-known short-range repulsion induced by vector-meson exchange. However, from the modern point of view such a mechanism looks doubtful.\(^8\) We will demonstrate below how the repulsive core effects can be reproduced within the framework of the quark and dibaryon models (see also Refs. \[26-31\]).

As is well known, the repulsive core effects are manifested very clearly in the channel \(^1S_0\), where the phase shifts become negative from collision energies \( E_{\text{lab}} \approx 250 \text{ MeV} \). The dibaryon model \[14, 15\] predicts for the \( S\)-wave \( N\bar{N} \) interaction the dominating six-quark configuration of the type \( |s^0p^2[42]_L^S\rangle \) with a two-quantum (\( 2\hbar \omega \)) excitation. When projecting this \( 2\hbar \omega \)-excited configuration onto the \( N\bar{N} \) channel, the \( N\bar{N} \) relative motion wavefunction \( \psi(r) \) automatically acquires an internal node. This node is rather stable and its position \( r_c \) moves only weakly with increase of collision energy (see Fig. 5). Moreover, the node position \( (r_c \approx 0.5 \text{ fm}) \) turns out to be very close to that of the traditional repulsive core.

In the dibaryon model \[14, 15\] appearance of a stationary node in the \( S\)-wave \( N\bar{N} \) interaction is provided by an orthogonality condition between the fully symmetric and the mixed-symmetry six-quark configurations, which is realized by adding a projecting operator to the \( N\bar{N} \) potential:

\[
V_{\text{rep}} = \lambda |\phi_0\rangle \langle \phi_0|,
\]

where \( \lambda \to \infty \), and \( |\phi_0\rangle \) is a projection of the fully symmetric wavefunction \( |s^0[6]| \) onto the \( N\bar{N} \) channel. The idea behind inclusion of this term into the \( N\bar{N} \) potential is that the six-quark configuration \( |s^0[6]\rangle \) is much smaller than the mixed-symmetry components. Hence, in case of \( S- \) and some of the \( P\)-wave channels \( (^3P_0, ^3P_1 \text{ and } ^1P_1) \), the effective \( N\bar{N} \) Hamiltonian \[8\] should be supplemented with the orthogonalizing pseudopotential \[10\] with a large positive coupling constant \( \lambda.9\)

![Figure 5: The nodal behavior of scattering functions in the \(^1S_0\) channel at different energies \( E_{\text{lab}} \): 10 MeV (solid curve), 100 MeV (dashed curve), 500 MeV (dash-dotted curve), and 1 GeV (dash-dot-dotted curve).](image)

Fortunately, this modification does not add to our model any free parameters, so that, the number of adjustable parameters for the bare dibaryon remains minimal. Thus, varying the position of only one complex pole in the \(^1S_0\) channel, one can achieve a very reasonable description of the phase shifts in a wide energy range. However, the resonance in this channel occurs to be very broad, hence, few different sets of parameters can be used to fit the phase shifts. In Fig. 5 we show the results for two possible positions of the bare dibaryon \( (M_D = 2.245 \text{ GeV and } 2.427 \text{ GeV}) \). Each of them allows to reproduce the real and imaginary parts of the \(^1S_0\) phase shifts nearly up to the respective bare resonance energy.

In particular, the potential parameters for the initial bare resonance with the lower mass \( M_D = 2.245 \text{ GeV} \) and the width\(^10\) \( \Gamma_0 = 0.1 \text{ GeV} \) have been taken as \( \lambda_1 = 0.51 \text{ GeV} \) and \( r_0 = 0.6 \text{ fm} \). For these potential parameters, the resulted \( S\)-matrix has two singularities corresponding to the well-known singlet deuteron just above the \( N\bar{N} \) threshold (at \( E = -0.065 \text{ MeV} \) on the unphysical complex energy sheet) and the high-lying resonance at \( M_{\text{th}} = 2.42 \text{ GeV} \) with a width of \( \Gamma_{\text{th}} = 0.44 \text{ GeV} \). We emphasize here once again that the mass and width of the bare \(^1S_0\) dibaryon were taken to be the same for de-

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\(^8\) See, e.g., Ref. \[28\] where the authors claim: “A literal attribution of the short-range repulsive core to vector meson exchange, as opposed to a phenomenological parametrization, of course involves a non sequitur: since the nucleons have radii \( \approx 0.8 \text{ fm} \) and the range of the vector exchange force is \( 1/m_{\omega,c} \approx 0.2 \text{ fm} \) one would have to superimpose the nucleon wavefunctions to reach the appropriate internucleon separations. The picture of distinct nucleons exchanging a physical \( \omega \)-meson at such a small separation is clearly a fiction...”

\(^9\) For the \( P\)-waves \( ^3P_0 \text{ and } ^3P_1 \) \( |\phi_0\rangle \) is a projection of the mixed-symmetry six-quark wavefunction \( |s^0p[51]_L^S\rangle \) onto the \( N\bar{N} \) channel.

\(^10\) Here the initial width parameters entering Eq. \[9\] were chosen as \( \alpha = 0.035 \text{ GeV} \) and \( \beta = 0.12 \text{ GeV} \).
These results give a strong indication for existence of a second high-lying dibaryon in the $^1S_0$ channel in addition to the near-threshold $^1S_0$ dibaryon (the singlet deuteron) predicted by Dyson and Xuong\cite{8} many years ago. It is interesting to note that while the second pole for the bare dibaryon has been postulated when constructing the effective potential (3), the first dibaryon pole appears in the course of dressing, i.e., when the coupling between the $NN$ and dibaryon channels gets switched on.

V. CONCLUSION

We have demonstrated in this work that the hypothesis about the dibaryon origin of the basic nuclear force makes it possible to describe properly the behaviour of both real and imaginary parts of $NN$ phase shifts (at least in some partial-wave channels characterized by a large inelasticity) in a rather wide energy range 0–600 MeV with only a few basic parameters for the initial “bare” dibaryon position. This result should be compared to the fact that the traditional so-called realistic $NN$ potentials as well as the modern EFT approach can describe only the real $NN$ phase shifts in the energy range 0–350 MeV.

What is even more interesting, the results of the present study also explain the strong channel dependence of $NN$ phase shifts, which must depend upon the dynamics of dibaryon states in the given partial-wave channel, their spin, isospin, parity, etc. In turn, the dibaryon dynamics is determined completely by QCD degrees of freedom. Without taking into account the intermediate dibaryons, the $NN$ interaction potential must have very complicated and non-transparent structure and adopt numerous terms with $S, L, L^2$, etc., operators to describe the empirical channel dependence of $NN$ phase shifts.

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