Abstract  Phase shifts and inelasticity parameters for $NN$ scattering in the partial-wave channels $^3S_1-^3D_1$ and $^1S_0$ at energies $T_{\text{lab}}$ from zero to about 1 GeV are described within a unified $NN$ potential model assuming the formation of isoscalar and isovector dibaryon resonances near the $NN^*(1440)$ threshold. Evidence for these near-threshold resonances is actually found in the recent WASA experiments on single- and double-pion production in $NN$ collisions. There, the excitation of the Roper resonance $N^*(1440)$ exhibits a structure in the energy dependence of the total cross section, which corresponds to the formation of dibaryon states with $I(J^\pi) = 0(1^+) + 1(0^+) + 1(0^+)$. These two $S$-wave dibaryon resonances may provide a new insight into the nature of the strong $NN$ interaction at low and intermediate energies.

Keywords  Nucleon-nucleon interaction · Dibaryon resonances · Single- and double-pion production · Roper resonance

1 Introduction

The traditional point of view on the strong $NN$ interaction at low energies ($T_{\text{lab}} \lesssim 350$ MeV) is based on the classic Yukawa concept \cite{1} suggesting $t$-channel meson exchanges between nucleons. Later on, this idea of Yukawa has been realized in the so-called realistic $NN$ potentials \cite{2,3,4}. Recently the realistic $NN$ potentials (of the second generation) have been replaced by the Effective Field Theory (EFT) which can treat single and multiple meson exchanges more consistently \cite{5,6}. However when the energy is rising beyond 350 MeV, the numerous inelastic processes enter the game and the application of the traditional approach meets many serious problems. Importantly, most of them are related to our poor understanding of the short-range $NN$ interaction and the corresponding short-range two- and many-nucleon correlations in nuclei and nuclear matter \cite{7}.

From the general point of view, these problems should be tightly interrelated to the quark structure of nucleons and mesons. On the other hand, the consistent treatment of the intermediate-energy $NN$ interaction, especially for inelastic processes, within the microscopic quark models is associated to so enormous difficulties \cite{8,9,10}, that nowadays we have to limit ourselves with some phenomenological or semi-phenomenological treatment. However it is still possible to use some hybrid approach and to combine the meson-exchange treatment for the long-range $NN$ interaction with the quark-motivated model for the intermediate- and short-range interaction \cite{11}. Such a model can be naturally based on the assumption about the six-quark bag (or dibaryon) formation at sufficiently short $NN$ distances, where the three-quark cores of two nucleons get overlapped with each other \cite{12}. Implementation of this idea does not require a detailed knowledge of the six-quark dynamics, but only needs the projection of the six-quark wavefunctions onto the $NN$ channel and an operator coupling the two channels of different nature, i.e., nucleon-nucleon and six-quark ones. So, in the $NN$ channel we can take into account only the peripheral meson-exchange interaction, while the influence of the internal $6q$ channel on the $NN$ interaction can be described by a simple mechanism of an intermediate dibaryon reso-
formance formation with appropriate $NN \leftrightarrow 6q$ transition form factors \cite{13,14,15}.

The dibaryon-induced mechanism for the short-range $NN$ interaction was initially suggested in Ref. \cite{12} and quite successfully applied to the description of $NN$ elastic scattering phase shifts and the deuteron properties in Refs. \cite{15,13}. However, these works did not consider the inelastic channels and also did not try to identify the $S$-matrix poles resulting from the fit of the phase shifts using the dibaryon resonances found experimentally. On the other hand, in recent years a number of dibaryon resonances have been discovered which are manifested most clearly in the inelastic processes \cite{16}.

In Refs. \cite{17,18,19}, we elaborated a unified model that can describe well both elastic phase shifts and inelasticities in $NN$ scattering in various partial-wave channels at laboratory energies from zero up to about 1 GeV. Thus, it has been shown that one can reproduce quite satisfactorily both elastic and inelastic $NN$ scattering phase shifts in a broad energy range using only a one-term separable potential with a pole-like energy dependence (with complex energy) for the main part of interaction and the one-pion exchange potential (OPEP) for the peripheral part of interaction. The model was applied to various partial-wave channels with $L > 0$: $^1D_2, ^3P_0, ^3P_2, ^3D_3-^3G_3$ and others, and the theoretical parameters (mass and width) of dibaryon resonances found from the fit of $NN$ scattering in these channels turned out to be very close to their experimental values \cite{17,18,19}.

However, the description of just $S$-wave $NN$ scattering (both elastic and inelastic) at low and intermediate energies should be especially sensitive to the assumptions made in the dibaryon-induced model. In fact, in the case of $S$-wave $NN$ scattering, absence of a centrifugal barrier allows the closest rapprochement of two nucleons to each other, so the short-range interaction in the $S$-wave channels provides the strongest impact to the phase shifts. This is especially true for inelastic scattering which, in turn, should be governed by the same mechanism as elastic scattering. Thus, it seems evident that the quark degrees of freedom should play a major role in the interaction mechanism. In the present paper, we study in detail just $S$-wave $NN$ scattering within the dibaryon-induced approach.

In fact, our model is based mainly on an assumption about the formation of dibaryon resonances which can be coupled to the various $NN$ channels. Therefore the existence (or nonexistence) of such states plays a decisive role in our approach. So, it is worth to briefly discuss the current experimental status of the dibaryon resonances before we can proceed further with the theoretical description of $NN$ scattering.

In recent years, many so-called exotic states have been observed in the charmed and beauty meson and baryon sectors. Common to these $X$, $Y$, $Z$ and pentaquark states is that they appear as narrow resonances near particle thresholds constituting weakly bound systems of presumably molecular character \cite{20}. A similar situation is also present in the dibaryonic sector, which can be investigated by elastic and inelastic $NN$ scattering.

Following the recent observation of the narrow dibaryon resonance $d^*(2380)$ with $I(J^P) = 0(3^+)$ in two-pion production \cite{21,22} and then in $NN$ elastic scattering \cite{23,24}, new measurements and investigations revealed and/or reconfirmed evidences for a number of states near the $N\Delta$ threshold. Among these the most pronounced resonance is the one with $I(J^P) = 1(2^+)$, mass $m \approx 2148$ MeV and width $\Gamma \approx 126$ MeV. Since its mass is close to the nominal $N\Delta$ threshold of 2.17 GeV and its width is compatible with that of the $\Delta$ itself, its nature has been heavily debated in the past, though its pole has been clearly identified in a combined analysis of $pp, \pi d$ scattering and $pp \leftrightarrow d\pi^+$ reaction \cite{25}. For a recent review about this issue see, e.g., Ref. \cite{16}. Very recently also evidence for a resonance with mirrored quantum numbers, i.e., $I(J^P) = 2(1^+)$ has been found having a mass $m = 2140(10)$ MeV and width $\Gamma = 110(10)$ MeV \cite{26,27}. Remarkably, both these states have been predicted already in 1964 by Dyson and Xuong \cite{28} providing agreement with the experimental findings both in mass and in width.

Whereas these two states represent weakly bound states relative to the nominal $N\Delta$ threshold and are of presumably molecular character with $N$ and $\Delta$ in relative $S$ wave, new evidence has been presented recently also for two states, where the two baryons are in relative $P$ wave: a state with $I(J^P) = 1(0^-)$, $m = 2201(5)$ MeV and $\Gamma = 91(12)$ MeV as well as a state with $I(J^P) = 1(2^-)$, $m = 2197(8)$ MeV and $\Gamma = 130(21)$ MeV \cite{29}. The values for the latter state agree with those obtained before in SAID partial-wave analyses \cite{25}. The masses of these $p$-wave resonances are slightly above the nominal $N\Delta$ threshold, which is understood as being due to the additional orbital motion \cite{30}. There is suggestive evidence for the existence of still further states like a $P$-wave $I(J^P) = 1(3^-)$ state, for which, however, the experimental situation is not yet as clear \cite{16}. It is also worth emphasising that the three resonances $1(2^+), 1(2^-)$ and $1(3^-)$ have been shown to give a sizeable contribution to the $pp \leftrightarrow d\pi^+$ cross sections and polarisation observables \cite{31} (the resonance
1(0−) is not allowed in this reaction by the parity and momentum conservation.

In the description of \( N N \) scattering within the dibaryon-induced model, we considered first the isovector partial channels \( 1^2D_2, 3^3P_2, 3^3P_3 \) and others, where the dibaryon resonances near the \( N \Delta \) threshold (respectively, \( 1(2^+) \), \( 1(2^−) \), \( 1(3^−) \), etc.) can be formed \cite{17,18}. We have shown that these resonances determine almost completely \( N N \) scattering in the respective partial channels at energies from zero to about 600–800 MeV (lab.). Then in the work \cite{19} \( N N \) scattering in the isoscalar \( 3^3D_3-3^3G_3 \) channels has been shown to be governed by the \( 0(3^+) \) dibaryon \( d^* (2380) \) which is located 80 MeV below the \( \Delta \Delta \) threshold (and thus can be treated not as a molecular-like but as a deeply bound \( \Delta \Delta \) state). By analogy, for the \( S \)-wave partial channels, with which we are concerned here, the respective dibaryons could be located near the \( NN^* (1440) \) threshold, since the Roper resonance \( N^* (1440) \) has the same quantum numbers as the nucleon, and an \( S \)-wave \( NN^* \) resonance can easily transform into an \( S \)-wave \( NN \) state. In comparison to \( N \Delta \) dibaryons which can couple to the isovector \( NN \) channels only, both isospin assignments \( I = 0 \) and \( I = 1 \) are allowed for the \( NN^* (1440) \) resonances. So, these resonances, if they exist, can couple to the \( 3^1S_1-3^1D_1 \) (the deuteron) and \( 1^3S_0 \) (the singlet deuteron) \( NN \) channels, respectively.

Fortunately, a strong indication of existence of these two dibaryon resonances near the threshold of the Roper resonance excitation have been found in the recent WASA experiments on single- and double-pion production in isoscalar and isovector \( NN \) collisions \cite{32,33}. It will be demonstrated below that the scenario of dibaryonic resonances near the \( N \Delta \) threshold is not unique, but is repeated at the \( NN^* (1440) \) threshold. And just these resonances determine the \( S \)-wave \( NN \) scattering at low and intermediate energies.

The paper is organised as follows. In Sec. 2 we briefly outline the theoretical formalism of the dibaryon-induced model for the \( NN \) interaction \cite{17,18,19} with some modifications necessary to apply it to \( S \)-wave scattering. Then in Sec. 3 we derive the dibaryon parameters fit to the phase shifts and inelasticities in the \( 3^1S_1-3^1D_1 \) and \( 1^3S_0 \) channels and compare them to the experimental data which are discussed in Sec. 4. We conclude in Sec. 5.

### 2 The dibaryon-induced model for the \( S \)-wave \( NN \) interaction

As is well known, the effective range approximation for the low-energy \( NN \)-scattering leads to the \( S \)-matrix poles near zero energy for the triplet \( 3^1S_1-3^1D_1 \) and singlet \( 1^3S_0 \) channels. According to the Wigner’s idea, one can treat these \( S \)-matrix poles as a result of an \( s \)-channel exchange by the deuteron or singlet deuteron (see Fig. [1]).

![Fig. 1](image.png)

**Fig. 1** Diagram illustrating the low-energy \( NN \) interaction due to \( s \)-channel exchange by the deuteron or singlet deuteron.

Then the question arises: whether such \( s \)-channel mechanism can provide description not only of low-energy but also intermediate-energy \( NN \) scattering? The answer is: surely, if instead of the deuteron pole in Fig. [1] one will imply a corresponding dibaryon pole at intermediate energy. The first attempt to treat the \( S \)-wave \( NN \) scattering at intermediate energies by the \( s \)-channel exchange by the dibaryon pole was undertaken at the beginning of 2000s within the framework of the dibaryon concept for the nuclear force \cite{13}. The peripheral meson-exchange \( NN \) interaction was described via the so-called external space (or channel) where one deals with nucleonic and mesonic degrees of freedom. The main short-range \( NN \) attraction is caused by a coupling between the external and internal channels, where the latter is treated by means of quark-glue (or string) degrees of freedom. The rigorous mathematical formalism to describe such quantum systems combining two Hilbert spaces (or channels) with completely different degrees of freedom was developed in the numerous papers of the Leningrad group \cite{34}. We refer the reader to Refs. \cite{13,14} where this approach was used to develop a dibaryon-induced \( NN \)-interaction model (referred to as a “dressed bag model”) based on a microscopic six-quark shell model in a combination with the well-known \( 3^1P_1 \) mechanism of pion production.

A deeper insight into the structure of the six-quark system in the internal channel may be gained from the quark-cluster picture \cite{35,36}, where two separated quark clusters, a tetraquark \( 4q \) and a diquark \( 2q \) are connected by a color string which can vibrate and rotate. In the quark shell-model language, such a state corresponds to the six-quark configuration \(|s^4p^2|^{12}_{2S}; L = 0, 2, ST \rangle \) with two quarks in the \( p \)-shell \cite{12,13}. Being transformed into the \( 4q-2q \) two-cluster state, it corresponds to the \( 2\omega \) excitation of the color string connecting two clusters. So, coupling between the external and internal channels corresponds to passing from
a bag-like $2\hbar\omega$-excited six-quark state to $NN$ loops in the external channel (see Fig. 2). Of course, the intermediate dibaryon can decay also into inelastic channels (other than $NN$). In our model, such decays are effectively taken into account through the width $\Gamma_D$ (see Eq. (2) below). So that, in Fig. 2 dibaryon decays into $NN^*$ channel are implicitly included in the dibaryon propagator as well.

Fig. 2 Graphical representation of the $NN$ scattering amplitude driven by the intermediate dibaryon $D$ formation in the $NN$ system.

In Refs. [17,18,19], the dibaryon-induced model has been generalised further to effectively include the inelastic processes. Below, for the readers’ convenience, we briefly outline the basic formalism of the dibaryon-induced model with a special emphasis on $S$-wave $NN$ scattering. As has been mentioned above, the total Hilbert space of the model includes the external and internal channels. The external channel corresponds to the relative motion of two nucleons, while the internal channel corresponds to the formation of the six-quark (or dibaryon) state. In the simplest case, the internal space is one-dimensional, and a single internal state $|\alpha\rangle$ is associated with the “bare dibaryon” having the complex propagator as well.

\[ H = \left( \begin{array}{c} h_{NN} \lambda|\Phi\rangle \langle \alpha | \\lambda|\alpha\rangle |\Phi\rangle \langle E_D|a\rangle \langle a| \end{array} \right) , \]  

where the transition form factor $|\Phi\rangle$ is defined in the external space and represents a projection of the total $6q$ wavefunction onto the $NN$ channel. In particular, in case of the coupled spin-triplet $NN'$ partial waves $^3S_1-^3D_1$, $|\Phi\rangle$ is a two-component column (see Ref. [19]).

The external Hamiltonian $h_{NN}$ includes the peripheral interaction of two nucleons which is given by the one-pion exchange potential $V_{OPEP}$. Here we use the same form and the same parameters of $V_{OPEP}$ as in Ref. [19]. For $S$-wave $NN$ scattering, one should also take into account the six-quark symmetry aspects leading to an additional repulsive term $V_{orth}$ in the $NN$ potential [18]. Thus, the external Hamiltonian is represented as a sum of three terms:

\[ h_{NN} = h_{NN}^0 + V_{OPEP} + V_{orth}, \]  

where $h_{NN}^0$ is the two-nucleon kinetic energy operator (which may include the Coulomb interaction for the $pp$ case) and $V_{orth}$ has a separable form

\[ V_{orth} = \lambda_0|\phi_0\rangle \langle \phi_0| . \]  

The symmetry-induced operator $V_{orth}$ was introduced for the first time in Ref. [27]. It corresponds to the full or partial exclusion of the space symmetric six-quark component $|s^0\rangle$ from the total $NN$ wavefunction and is needed to fulfill the orthogonality condition between the small $|s^6\rangle$ and the dominating mixed-symmetry $|s^4p^2\rangle$ components in the $NN$ system. It has been shown [38] that the operator $V_{orth}$ plays the role of the traditional $NN$ repulsive core. In fact, this $s^6$-eliminating potential provides a stationary node in the $NN$ wavefunctions at different energies, and the position of the node corresponds to the radius of the repulsive core [39].

To satisfy the orthogonality condition strictly, one has to take the limit $\lambda_0 \to \infty$ in $V_{orth}$. However, since the $S$-wave $NN$ channels have a strong coupling to the $S$-wave $NN^*(1440)$ channels near the Roper resonance excitation threshold, the $2\hbar\omega$ excitation in $NN$ relative motion can pass into the $2\hbar\omega$ inner monopole excitation of the Roper resonance $N^*(1440)$. Thus, for such a strong coupling, there should not be a strict orthogonality condition for the symmetric configuration $|s^0\rangle$ at energies near the resonance, and the value of $\lambda_0$ should be finite. There is another good reasoning to this point. The $S$-wave dibaryon state located near the $NN^*(1440)$ threshold can decay into both $NN$ and $NN^*$ channels. While the relative-motion wavefunction in the $NN$ channel has a stationary node at $r_c = 0.5$ fm similarly to the low-energy $NN$ scattering the $NN^*$ wave function has not got a node because the $2\hbar\omega$ excitation in the initial six-quark wave function passes into $2\hbar\omega$ inner excitation in the Roper state itself. Hence, for the channel $NN^*$, the projection operator $V_{orth}$ is not needed. The especially strong mixing of the $NN$ and $NN^*$ channels happens just in the near-threshold area where the effect of $V_{orth}$ almost disappears.

So that, we use here the orthogonalising term $V_{orth}$ with the finite values of $\lambda_0$. It provides a node in the $NN$ relative motion wavefunctions at small energies, but at intermediate energies, $NN$ scattering states have some admixture of the nodeless state $|\phi_0\rangle$. So, in this approach, the finite properly chosen value of $\lambda_0$ provides

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Footnotes:

1. For the coupled spin-triplet channels $^3S_1-^3D_1$, we use a bit lower cutoff parameter $A_{NN} = 0.62$ GeV (instead of 0.65 GeV employed in Ref. [19]) which allows for a better fit of the $^3D_1$ phase shift.

2. In the quark shell-model language, the $N^*(1440)$ structure corresponds to the mixture of the $3q$ configurations $(0s)^2 - (1p)^2$ and $(0s)^2 - 2s$, both carrying $2\hbar\omega$ excitation.

3. In particular, a resonance state in the $NN$ channel may have a noticeable overlap with the state $|\phi_0\rangle$. The detailed study of this formalism will be published elsewhere.
an effective account of the strong coupling between the \( NN \) and \( NN^*(1440) \) channels.

After excluding the internal channel, one gets the effective Hamiltonian in which the main attraction is given by the energy-dependent pole-like interaction:

\[
H_{\text{eff}}(E) = h_{NN} + \frac{\lambda^2}{E - E_D} |\Psi\rangle \langle \Phi|.
\]

By using the separable form for the energy-dependent part of interaction, one can find explicitly an equation for the poles of the total \( S \)-matrix (see details in Refs. [18, 19]):

\[
Z - E_D - J(Z) = 0,
\]

where the function \( J(Z) \) is determined from the matrix element of the external Hamiltonian resolvent \( g_{NN}(Z) = [Z - h_{NN}]^{-1} \), i.e., \( J(Z) = \lambda^2 |\Phi| g_{NN}(Z) |\Psi\rangle \).

Finally, for the effective account of inelastic processes, we introduce the imaginary part of the internal pole position \( E_D = E_0 - i \Gamma_D/2 \), which is energy-dependent and describes the possible decays of the “bare” dibaryon into all inelastic channels (i.e., except for the \( NN \) one). For a single decay channel, the width \( \Gamma_D \) can be represented as follows:

\[
\Gamma_D(\sqrt{s}) = \begin{cases} 0, & \sqrt{s} \leq E_{\text{thr}}; \\ \Gamma_0, & \sqrt{s} > E_{\text{thr}} \end{cases},
\]

where \( \sqrt{s} \) is the total invariant energy of the decaying resonance, \( M_0 \) is the bare dibaryon mass, \( E_{\text{thr}} \) is the threshold energy, and \( \Gamma_0 \) defines the partial decay width at \( \sqrt{s} = M_0 \). For the \( S \)-wave dibaryon resonance located near the \( NN^*(1440) \) threshold, the dominant decay channel is \( D \to NN^*(1440) \). Here we take into account only the main decay mode of the Roper resonance \( NN^*(1440) \to \pi NN \), and thus the main decay channel of the dibaryon is \( D \to \pi NN \). For such a case, the parametrization of the function \( F(\sqrt{s}) \) in Eq. (6) has been introduced in Ref. [18]:

\[
F(\sqrt{s}) = \frac{1}{s} \int_{2m}^{\sqrt{s}-m} dM_{NN} q^{2L+1} k^{2L+1} (q^2 + A^2)^{L+1} (k^2 + A^2)^L s^{L+1}.
\]

Here \( q = \sqrt{(s - m_N^2 - M_{NN}^2)^2 - 4m_N^2 M_{NN}^2 / 2\sqrt{s}} \) and \( k = 1/2 \sqrt{M_{NN}^2 - 4m_N^2} \) are the pion momentum in the total center-of-mass frame and the momentum of the nucleon in the center-of-mass frame of the final \( NN \) subsystem with the invariant mass \( M_{NN} \), respectively. In Eq. (7), the high momentum cutoff parameter \( A \) is used to prevent an unphysical growth of the width \( \Gamma_{\text{inel}} \) at high energies. The possible values of the pion orbital angular momentum \( l_\pi \) with respect to the \( NN \) subsystem and the orbital angular momentum of two nucleons \( L_{NN} \) are restricted by the total angular momentum and parity conservation. Below, the particular values of \( l_\pi, L_{NN} \) and \( A \) are adjusted to get the best fit of inelasticity parameters in the partial \( NN \) channels in question.

### 3 Results for the \( ^3S_1 - ^3D_1 \) and \( ^1S_0 \) partial-wave channels

In this section, we present the results of calculations for the \( NN \) scattering phase shifts and inelasticity parameters in the lowest partial-wave channels, viz. \( ^3S_1 - ^3D_1 \) and \( ^1S_0 \), within the dibaryon-induced model.

For the model form factors entering Eqs. (1) and (3), we have employed the harmonic oscillator functions with the orbital momentum \( L \) and the radial quantum number \( n \) equal to the number of nodes in the \( NN \) relative motion wavefunction. In particular, \( |\Phi_0\rangle \) has the form of the \( S \)-wave state with \( n = 0 \) and an effective range \( r_0 \). The “dibaryon” form factor \( |\Phi\rangle \) has two components corresponding to \( S \) and \( D \) waves, i.e.,

\[
|\Phi\rangle = \left( \frac{\alpha}{|\phi_S\rangle} + \frac{\beta}{|\phi_D\rangle} \right),
\]

where \( \alpha^2 + \beta^2 = 1 \). Here \( |\phi_S\rangle \) has the same effective range \( r_0 \) as \( |\phi_0\rangle \), but \( n = 1 \), so, these two functions are orthogonal to each other. The \( D \)-wave part of \( |\Phi\rangle \) is a nodeless function with the effective range \( r_D \). The potential parameters used for both spin-triplet and spin-singlet partial-wave channels are listed in Tab. [1] Here \( \lambda_S \equiv \alpha \lambda \) and \( \lambda_D \equiv \beta \lambda \) for the spin-triplet channel. For the dibaryon width defined by Eqs. (6) and (7), we used the values \( l_\pi = 0, L_{NN} = 1 \) and \( A = 0.6 \) GeV/c for the \( ^1S_0 \) channel and \( l_\pi = 1, L_{NN} = 2 \) and \( A = 1.8 \) GeV/c for the coupled \( ^3S_1 - ^3D_1 \) channels. The concrete values of these parameters are important mainly for the fit of inelasticities in the near-threshold region and have a little impact on the overall fit quality.

**Table 1** Parameters of the dibaryon model potential for the lowest spin-triplet and spin-singlet \( NN \) partial-wave channels.

| Channel | \( \lambda_0 \) MeV | \( r_0 \) fm | \( \lambda_S \) MeV | \( \lambda_D \) MeV | \( r_D \) fm | \( M_0 \) MeV | \( \Gamma_0 \) MeV |
|---------|-------------------|------------|----------------|----------------|-----------|-----------|-------------|
| \(^3S_1 - ^3D_1\) | 165 | 0.475 | 248.1 | 65.9 | 0.6 | 2275 | 80 |
| \(^1S_0\) | 165 | 0.48 | 274.2 | - | - | 2300 | 40 |

The partial phase shifts and mixing angle for the coupled channels \(^3S_1 - ^3D_1\) are shown in Fig. [3] in comparison with the single-energy (SE) solution of the SAID partial-wave analysis (PWA) [10]. It is seen from Fig. [3] that within the dibaryon model we can reproduce the PWA data on the \(^3S_1 - ^3D_1\) partial phase shifts and mix-
The partial phase shifts for the spin-singlet channel $^{1}S_0$ calculated with the model parameters which are rather close to those used for the spin-triplet case (see Tab. 1) are shown in Fig. 4 in comparison with the SAID SE data. Again the dibaryon model allows for the very good description of the partial phase shifts at energies from zero up to about 1.2 GeV.

The partial phase shifts for the spin-singlet channel $^{1}S_0$ found within the dibaryon model (solid curves) in comparison with the single-energy SAID PWA [40] (filled circles) and results for the pure OPEP (dash-dotted curves). The comparison of inelasticities for the $NN$ channel $^{1}S_0$ calculated with the model parameters which are rather close to those used for the spin-triplet case (see Tab. 1) are shown in Fig. 4 in comparison with the SAID SE data. Again the dibaryon model allows for the very good description of the partial phase shifts at energies from zero up to about 1.2 GeV.

The comparison of inelasticities for the $S$-wave channels with the SAID single-energy data is presented in Fig. 3 (a) and (b). Here we see reasonable agreement for the $S$-wave inelasticity parameters with the PWA data up to the energies corresponding to the resonance position ($T_{th} ≃ 0.9$ GeV). Thus, the same single-pole model of interaction can reproduce almost quantitatively both elastic and inelastic $NN$ scattering in $S$ waves in a broad energy range from zero up to about 1 GeV.

In Figs. 3-5 the contribution of the pure OPEP is shown by dash-dotted curves. It is clearly seen that just the dibaryon excitation mechanism allows for a reasonable description of both partial phase shifts and inelasticities for $S$-wave $NN$ scattering. The coupling with a dibaryon in the $D$-wave component of the spin-triplet channel $^{3}S_1^{*}D_1$ is weaker, so the dibaryon mechanism makes some important contribution here only above the inelastic threshold (see Fig. 3 (b)). The situation here is very similar to that for the $^{3}D_3^{*}G_3$ partial-wave channels studied in Ref. [19].

It is extremely interesting that the $S$-matrices for the model $NN$ potentials in both singlet and triplet partial channels have two poles. For the $^{3}S_1^{*}D_1$ case, the first pole corresponds to the bound state, i.e., the deuteron, which is reproduced rather accurately. The second pole here corresponds to the dibaryon resonance with the parameters:

$$M_{th}(^{3}S_1^{*}D_1) = 2310 \text{ MeV, } \Gamma_{th}(^{3}S_1^{*}D_1) = 157 \text{ MeV}.$$  \hspace{1cm} (8)

For the $^{1}S_0$ channel, the first pole is the well-known singlet deuteron state, while the position of the second one is:

$$M_{th}(^{1}S_0) = 2330 \text{ MeV, } \Gamma_{th}(^{1}S_0) = 51 \text{ MeV}.$$  \hspace{1cm} (9)

Both these resonance positions are rather close to the $NN^{*}(1440)$ threshold. As will be shown below, the resonance parameters given in Eqs. 8 and 9 also turn out to be very close to the values derived from the recent single- and double-pion production experiments (see Sec. 4). However, the inaccuracy in description of inelasticity parameters at energies above the resonance position in the considered $NN$ partial-wave channels as well as a too narrow width of the dibaryon resonance in

\[^{4}\text{The difference between the resonance parameters found here for the }^{1}S_0\text{ channel from the preliminary ones obtained in Ref. [13] is due to the use of the finite } \lambda_0 \text{ in the orthogonalising potential } V_{orth}.\]
the $^1S_0$ channel show that a more detailed treatment of inelastic processes is required within the dibaryon model.

As $s$-channel resonances, the two predicted dibaryon states have to display a counter clockwise looping in the Argand diagrams of amplitudes in $^1S_0$ and $^3S_1$ partial waves. For these partial waves, two $S$-wave trajectories are shown by the solid lines in Fig. 5 in comparison with the different SAID PWA solutions [40]. In fact, we observe the counter clockwise loopings for the amplitudes found within the dibaryon model indicating the resonance presence in both cases.

Since these resonances are highly inelastic, the resonance loops are rather tiny. The theoretical predictions should be compared with three PWA solutions of the SAID group, viz., the single-energy as well as the global solutions AD14 and SM16 [40]. The scatter within the single-energy data as well as the differences among the various SAID solutions may serve as an indication for inherent ambiguities in the partial-wave analysis, especially for the $^3S_1$ channel. In fact, the differences between two recent solutions SM16 and AD14 are of the same order as those differences between theoretical loops and each of the above SAID solutions. Hence, the absence of the loops in the current SAID solutions cannot argue against the suggested dibaryon resonances.

We note that the situation here is much different from that for the dibaryon resonance $d^*(2380)$. In case of the latter, the partial waves $^3D_3$ and $^3G_3$ were involved, which both carry large orbital angular momentum and hence have a large impact on the analyzing power. Since this observable is the only one, which solely consists of interference terms, it is predestinated to exhibit substantial effects even from tiny resonance admixtures in partial waves. Unfortunately, we deal here with $S$-wave resonances, which make no contribution to the analyzing power due to the missing orbital angular momentum. Hence this key observable for revealing loops and resonances is not working here. The only way out of this dilemma is to look into reactions, where these highly inelastic resonances decay to, namely single- and double-pion production. We discuss these processes in the next Section.

4 The Roper excitation in $NN$ induced single- and double-pion production and the near-threshold dibaryon resonances

The Roper resonance $N^*(1440)$ excitation appears usually quite hidden in the observables and in most cases can be extracted from the data only by sophisticated analysis tools like partial-wave decomposition. By contrast, it can be observed free of background in $NN$-induced isoscalar single-pion production, where the overwhelming isovector $\Delta$ excitation is filtered out by isospin selection as demonstrated by recent WASA-at-COSY results [42] for the $NN \rightarrow [NN\pi]_{I=0}$ reaction. Though the primary aim of this experiment was the search for a decay $d^*(2380) \rightarrow [NN\pi]_{I=0}$, it also covers the region of the Roper excitation, which is discussed here.

Since the $\Delta$ excitation is filtered out by the isospin condition, there is only a single pronounced structure left in the isoscalar nucleon-pion invariant mass spectrum as seen in Fig. 6 of Ref. [42], which peaks at $m \approx 1370$ MeV revealing a width of $\Gamma \approx 150$ MeV. These values are compatible with the pole values for the Roper resonance deduced in diverse $\pi N$ and $\gamma N$ studies [41]. Our values for the Roper peak also are in good agreement with earlier findings from hadronic $J/\Psi \rightarrow NN\pi$ decay [42] and $\alpha N$ scattering [43].
The energy excitation function of the measured \(NN\)-induced isoscalar single-pion production cross section is displayed in Fig. 7. Near threshold the Roper resonance is produced in \(S\) wave in relation to the other nucleon, whereas the pion from the Roper decay is emitted in relative \(p\) wave. Hence we expect for the energy dependence of the total cross section in the isoscalar \(NN\rightarrow[NN\pi]_{I=0}\) channel a threshold behavior like for pion \(p\) waves — as is actually born out by the explicit calculations for the \(t\)-channel Roper excitation in the framework of the modified Valencia model [32,46]. These calculations are displayed in Fig. 7 by the dashed line, which is arbitrarily adjusted in height to the data. The data [32–45] presented in Fig. 7 follow this expectation by exhibiting an increasing cross section with increasing energy up to about \(\sqrt{s} \approx 2.30\) GeV. Beyond that, however, the data fall in cross section in sharp contrast to the expectation for a \(t\)-channel production process. The observed behavior is rather in agreement with a \(s\)-channel resonance process as expected for the formation of a dibaryonic state at the \(NN^*\) threshold. Due to the relative \(S\) wave between \(N\) and \(N^*\) as well as due to the isoscalar nature of this system, it must have the unique quantum numbers \(I(J^P) = 0(1^+)\). From a fit of a simple Lorentzian to the data we obtain \(m = 2315(10)\) MeV and \(\Gamma = 150(30)\) MeV. The large uncertainty on the latter results from the large uncertainties of the data at lower energies (for a fit with a Gaussian, which leads to a width of 170 MeV, see Ref. [32]). For a more detailed treatment of the resonance structure one would need to use a momentum-dependent width, which takes into account the nearby pion production threshold and lowers the resonance cross section at the low-energy side.

A very similar situation is also observed in \(NN\)-induced two-pion production. The situation is particularly clear in the \(pp\rightarrow pp\pi^0\pi^0\) reaction, the total cross section of which is plotted in Fig. 8. Since ordinary single-\(\Delta\) excitation is excluded here due to the necessary production of two pions, the Roper excitation is the only resonance process at low energies. Hence we would again expect a phase-space-like (dotted line) growth of the cross section, which is also born out by detailed model calculations (dash-dotted line) [16,32]. But the data follow this trend up to \(T_p \approx 0.9\) GeV (\(\sqrt{s} \approx 2.3\) GeV). Then the data level off before they increase again, when the next higher-energetic process, the \(t\)-channel \(\Delta\Delta\) process with double-\(p\)-wave emission (dashed line) starts. Isospin decomposition of the data.
in the various \(NN\pi\) channels tells us that the energy dependence of the Roper excitation is experimentally given by the filled star symbols in Fig. 3 [33]. Again we see a resonance-like energy dependence, which indicates a \(NN^*(1440)\) molecular system also in this case, but now with quantum numbers \(I(J^P) = 1(0^+)\), \(m \approx 2320\) MeV and \(\Gamma \approx 150\) MeV. We note that the fading away of the Roper excitation at energies beyond \(T_p \approx 0.9\) GeV is also in agreement with the analysis of the corresponding differential cross sections [51,53].

![Fig. 8](image)

**Fig. 8** (Color online) Energy dependence of the total \(pp \rightarrow ppp^\pi^0\) cross section. Shown are the data from CELSIUS/WASA [78] as well as WASA-at-COSY [31] (filled circles), PROMICE/WASA [29] (filled squares), and earlier work [19,20] (open symbols). The dotted and dashed-dotted lines show the expected energy dependence of simple phase-space and modelled Roper excitation [65], respectively. The dashed line shows the \(t\)-channel \(\Delta\Delta\) excitation [40,51], whereas the filled stars display the result of the isospin decomposition for \(N^*\) excitations [33]. Here, the first structure is due to the Roper \(N^*(1440)\) excitation. The rise at higher energies signals higher-lying \(N^*\) excitations.

5 Conclusions

We have shown within the dibaryon-induced model for \(NN\) scattering that the \(NN\) interaction in the basic spin-singlet and spin-triplet \(S\)-wave partial channels at energies \(T_{lab}\) from zero up to about 1 GeV is governed by the formation of the \(I(J^P) = 0(1^+)\) and \(1(0^+)\) dibaryon resonances near the \(NN^*(1440)\) threshold. This work continues a series of the previously published papers [17,18,19] where the \(NN\) interaction in higher partial waves was shown to be dominated by the intermediate dibaryon excitation (supplemented by the peripheral one-pion exchange) in the respective partial channels near the \(N\Delta\) or \(\Delta\Delta\) thresholds.

From the energy dependence of \(NN\)-induced isoscalar single-pion and isovector double-pion production we see also that both isospin-spin combinations in the \(NN^*(1440)\) system lead obviously to dibaryonic threshold states at the Roper excitation threshold — analogous to the situation at the \(\Delta\) threshold. However, compared to the situation there the Roper excitation cross sections discussed here are small. Since these structures decay mainly into inelastic channels, their partial decay width into the elastic \((NN)\) channel should be only a small fraction of the total width, similarly to the respective branching ratio for the \(N\Delta\) near-threshold states [16]. Despite this fact, our results show that the contributions of these dibaryon states to the low- and intermediate-energy \(NN\) elastic scattering are dominating.

On the other hand, at energies below the inelastic thresholds \(NN\pi\) and \(NN\pi\pi\), decay of the dibaryons into these channels is forbidden. However, a strong coupling between the \(NN\) and the closed (virtual) channels like \(N\Delta\), \(NN^*(1440)\), \(NN\pi\) and \(NN\pi\pi\) is still possible. So, at low energies \((T_{lab} \lesssim 350\) MeV\) the coupling of the dibaryon to these closed channels appears to be strong and thus the whole picture of the \(NN\) interaction at these energies is dominated just by this coupling. This explains how the intermediate dibaryon formation near the nucleonic resonance threshold can be the leading mechanism for the \(NN\) interaction at low energies. When the collision energy is rising and the inelastic channels open, the same intermediate dibaryons provide single- and double-pion production. Thus, in the dibaryon-induced approach to the \(NN\) interaction, the elastic and inelastic \(NN\) collision processes have a common origin and can be described via a common mechanism. These results may provide a novel insight into the nature of the \(NN\) interaction at low and intermediate energies and should be confirmed by further experimental and theoretical research.

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