Pionic Content of $\rho NN$ and $\rho N\Delta$ Vertex Functions

Q. Haider
Physics Department, Fordham University, Bronx, NY 10458

L. C. Liu
Theoretical Division, T-2, Los Alamos National Laboratory
Los Alamos, NM 87545

Abstract

The dynamical content of $\rho NN$ and $\rho N\Delta$ vertex functions is studied with a mesonic model. A set of coupled integral equations satisfied by these vertex functions were solved self-consistently. These solutions indicate that the dominant mesonic content arises from di-pion dynamics. With the experimentally determined pion-baryon-baryon coupling constants and ranges as input, the model predicts a $g_{\rho NN}$ that agrees with the meson-exchange-potential results. On the other hand, it predicts a smaller $f_{\rho N\Delta}$ and much softer form factors. Implications of the findings on the use of phenomenological coupling constants in nuclear reaction studies are discussed.

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Meson-exchange theories have been very successful in describing nuclear phenomena in the region of low and intermediate momentum transfers, where practical methods for solving non-perturbative QCD are yet to be fully developed. Within the framework of the meson-exchange theory, understanding the dynamics associated with the $\rho NN$ and $\rho N\Delta$ vertices is a subject of great interest. It is well known that $\rho$-exchange gives rise to a nucleon-nucleon ($NN$) tensor force which has a sign opposite to that arising from the exchange of $\pi$. The partial cancellation between these tensor forces makes it indispensable to include $\rho$ in the hadronic description of nuclei and nuclear matter. The important role of $\rho$ in nuclear structure has further motivated the study of its role in nuclear reactions. Studying the $\rho$ is, however, complicated by the fact that, unlike $\pi N \rightarrow \Delta$, the $\rho N \rightarrow N$ and $\rho N \rightarrow \Delta$ processes do not occur in free space. Consequently, information on these reactions were mainly obtained from fitting $NN$ phase shifts with the use of meson-exchange potentials (MEP). It is, therefore, quite likely that the MEP parameters may have to be modified for studying processes other than $NN$ scattering. Indeed, analyses of deep inelastic scattering data indicate that the range of a pion monopole form factor cannot exceed 650 MeV.[1][2] This upper limit is nearly a factor of two smaller than the MEP ranges. Smaller ranges have also been predicted as a necessary outcome of the boson nature of the pion.[3] A recent study of the $\rho(p, n)\Delta^{++}$ reaction has shown that the inclusion of $\rho$-exchange mechanisms actually worsens the fit.[4] This disagreement raises the interesting question as to what are the appropriate $\rho$-meson parameters to be used in meson production. These developments have motivated us to carry out a detailed analysis of the $\rho NN$ and $\rho N\Delta$ vertices by means of a dynamical model that does not treat the rho-baryon-baryon ($\rho BB'$) processes as contact interactions. The main goal of our study is to identify important dynamical contents of the phenomenological coupling constants.

Our model for the $\rho NN$ vertex is illustrated in fig.1, where the wavy, dashed, thin, and thick lines denote, respectively, the $\rho$, $\pi$, $N$, and $\Delta$. The initial and final nucleons are labelled $a$ and $c$, while the intermediate baryons are labelled $b$ and $b'$. Figure 1 corresponds to the equation

$$F_{\rho a;c} = F_{\rho;\pi\pi}G^{(+)\pi_1}_1G^{(+)\pi_2}_2 F_{\pi_1;\alpha;b}G^{(+)b}_b F_{\pi_2;\beta;c}' + F_{\pi;\pi\pi}G^{(+)\pi_1}_1G^{(+)\pi_2}_2 F_{\rho;b;b'}G_{b'}^{(+)b'} F_{\pi;\beta;c} \equiv B_{\rho a;c} + F'_{\rho a;c} ; \quad (1)$$
where $F_{\rho ac}$ denotes the $\rho a \to c$ vertex function. A similar notation applies to the other vertex functions, with the indices $b, b'$ representing either an $N$ or a $\Delta$. In Eq.(1) $B_{\rho ac}$ denotes the Born amplitude given by the two triangle diagrams in fig.1, while $F'_{\rho ac}$ corresponds to pion corrections of the $\rho b \to b'$ vertices, as shown by the four $\pi$-loop diagrams. Since the solutions of the model, $F_{\rho ij}$ $(i, j = a, b, b', c)$, appear on both sides of Eq.(1), we must solve a set of four coupled equations of the form of Eq.(1) with the baryon indices $ac$ equal to $NN$, $N\Delta$, $\Delta N$, and $\Delta\Delta$, respectively. For succinctness, only the equation corresponding to $ac = NN$ is illustrated in fig.1. We have solved the coupled equations by iteration, and a convergence was obtained after five iterations. In principle, one should also include the $\rho$-loop corrections of the vertices. Because $M_\rho > M_\pi$, these corrections will be much less important than the $\pi$-loop corrections which, as we shall see, are quite small. They were thus neglected in our calculations.

It is worth noting that diagrams similar to ours have recently been considered in ref.[5]. We emphasize, however, that there exist important differences between the analysis in ref.[5] and ours. In ref.[5] the leading contributor was assumed to be the phenomenological three-branch $\rho NN$ vertex while the triangle and loop diagrams were considered as corrections and treated perturbatively. Moreover, phenomenological MEP $\rho BB'$ coupling constants were used as inputs and no coupled equations were solved. As such, it is equivalent to evaluating only first-order corrections to the MEP $\rho NN$ coupling constant. Consequently, the calculation does not address the dynamics leading to the phenomenological coupling constants. In contrast, the analysis presented in this paper considers the triangle diagrams as the driving terms and the loop diagrams as corrections. By solving the coupled equations, the $\rho BB'$ coupling constants and ranges are obtained as the self-consistent solutions of the model considered. Hence, these solutions contain valuable information on the dynamical content of the coupling constants.

We used the following partial-wave decomposition for the $\rho BB'$ vertex functions:

$$F_{\rho ac}(p_\rho, p_a; p_c) = < I_a t_a, 1t_\rho | (I_a 1)_L d_c > \sum_{S_m L_m} < J_a \nu_a, 1\nu_\rho | (J_a 1)_S m_S >$$

$$\times < S m_S, L m_L | (SL) J_c \nu_c > Y^*_L m_L (\hat{p}) F(p_\rho, p_a, p_c) ,$$  (2)
with the radial vertex function $\mathcal{F}$ parametrized by

$$
\mathcal{F}(p_\rho, p_a, p_c) = \mathcal{G}(L)_{pa;c} \frac{1}{\sqrt{2w_c}} p^T v_L (\Lambda_{pa;c}, p).
$$

(3)

The $p_i$ ($i = \rho, a, c$) denote the four-momenta of the $i$th particle, and $p_i$ is its spatial part; $p$ denotes the magnitude of the relative momentum between $\rho$ and $a$, and $w_c$ is the invariant mass of $c$. The two Clebsch-Gordan coefficients specify, respectively, the isospin and angular momentum coupling schemes with $I$ and $J$ denoting the isospin and spin of the particles, and $t$, $\nu$ their $z$-components. The vertex function $\mathcal{F}$ is a scalar and depends on two independent four-momenta which can be chosen as $p_\rho + p_a$ and $p_\rho - p_a$. It is convenient to work in the $\rho - a$ c.m. system where the external four-momenta have the simple expressions $p_\rho = (\rho^0, p)$, $p_a = (a^0, -p)$, and $p_c = (w_c, 0)$. Clearly, $w_c = a^0 + \rho^0$. From the two independent four-momenta one can form three independent invariant scalars. We will put the baryon $a$ on the mass shell so that $a^0 = E_a(p) = \sqrt{m_a^2 + p^2}$. Consequently, $\mathcal{F}$ will depend only on two independent variables which we choose as $w_c (\equiv w)$ and $p (\equiv |p|)$. Because the parametrization used in Eq.(3) is not the most general one, we should expect $\mathcal{G}$ and $\Lambda$ to depend on $w$. This $w$-dependence has important physical consequences and will be discussed later. In the literature, vertex functions have often been made to depend only on one variable. We emphasize that the one-variable dependence is exact only when two of the three particles are on their mass shells.

Parity conservation limits the values of orbital angular momentum in Eq.(2) to $L = 1$ for $\rho N \rightarrow N$ and $\rho N \rightarrow \Delta$, and to $L = 1$ and $3$ for $\rho \Delta \rightarrow N$ and $\rho \Delta \rightarrow \Delta$. As the Lagrangian models in the literature consider only $L = 1$, we shall solve the coupled equations in the $p$-wave channel and omit, henceforth, the index $L$ of $\mathcal{G}$ and $v$. Parity conservation also limits the $\rho \rightarrow \pi \pi$, $\pi N \rightarrow N$, $\pi N \rightarrow \Delta$, and $\pi \Delta \rightarrow N$ processes to relative $p$-wave interactions alone. On the other hand, $\pi \Delta \rightarrow \Delta$ can have both $p$- and $f$-wave interactions. Again, only the $p$-wave interactions will be retained in the calculations in order to make a close connection to the Lagrangian models. The $\rho \pi \pi$ vertex function is parametrized as

$$
F_{\rho \pi \pi}(p_\rho, p_1, p_2) = \sum_{t_1 t_2} C_I \sum_m \frac{G_{\rho \pi \pi}}{2M_\rho} \bar{h}(\Lambda_{\rho \pi \pi}, k) Y^{*}_{1m}(\hat{k}) \sqrt{2\omega_\pi(p_1)2\omega_\pi(p_2)},
$$

(4)

where $C_I \equiv <1t_1, 1t_2 | (11)t_\rho >, \bar{h} \equiv kr/(1 + k^2 r^2)^2$ with $r \equiv 1/\Lambda_{\rho \pi \pi}$, and $k = |k|$ is the
Upon introducing Eqs. (3)-(5) into Eq. (1) and projecting out the $\pi\pi$ relative momentum. The $\pi_1ab$ vertex function is parametrized by

$$F_{\pi_1ab}(p_1, p_a; p_b) = \sum_{I_a I_b J_{b1} m_1} C_I C_J \frac{G_{\pi a b}}{2\sqrt{\omega_\pi}(p_1)} h(\Lambda_{\pi a b}, \kappa) Y_{1m_1}^*(k),$$

where $C_I \equiv< I_a | t_a, 1t_1 | (J_a 1) I_b | t_b >$, $C_J \equiv< J_a | \nu_a, 1m_1 | (J_a 1) t_b | \nu_b >$, $\kappa$ is the $\pi a$ relative momentum, $m_1$ is the z-component of the corresponding orbital angular momentum, and $h$ is a dipole form factor defined by $h \equiv (\kappa/M_\pi) [(\Lambda^2 - M_\pi^2)/(\Lambda^2 + \kappa^2)]^2$. The vertex functions $F_{\pi_2 b c}, F_{\rho b b^\prime}, F_{a: \pi b},$ and $F_{\pi b^\prime: c}$ are parametrized in a form similar to Eq. (3). The triangle as well as the loop diagrams depend on one four-momentum integration variable $q = (q^0, \mathbf{q})$. By closing the contour of $q^0$ along a semi-circle in either the upper or the lower half-plane, we can carry out the $q^0$ integration analytically. Upon introducing Eqs. (3)-(5) into Eq. (1) and projecting out the $(L = 1)$ angular momentum dependence, we obtain a set of four coupled equations for $v$.

One should note that the value of a coupling constant depends on the parametrization convention of a theory. It is, therefore, useful to relate the $G$’s of the present partial-wave formalism to the coupling constants of other works. To this end, we have used the S-matrix convention of ref.[9]. In relation to the Lagrangian model and the parametrization employed in ref.[6], we obtain $G_{\pi NN} = \sqrt{(3/2\pi^2)} f_{\pi NN}$ and $G_{\pi N\Delta} = f_{\pi N\Delta}/\sqrt{6\pi^2}$. For the $\pi\Delta\Delta$ vertex, $G_{\pi\Delta\Delta} = (5/2) f_{\pi\Delta\Delta}'/\sqrt{6\pi^2}$ where $f_{\pi\Delta\Delta}' \equiv (M_\pi/2M_\Delta) g_{\pi\Delta\Delta}$ and $g_{\pi\Delta\Delta}$ is the axial-vector coupling constant defined in ref.[8]. With respect to the parametrization in ref.[6], $G_{\rho\pi\pi}/\sqrt{2M_\rho} = g_{\rho\pi\pi}$ and $G_{\pi\Delta\Delta} = (15/4) f_{\pi\Delta\Delta}/\sqrt{6\pi^2}$. Hence, $f_{\pi\Delta\Delta} = (2/3) f_{\pi\Delta\Delta}'$. In addition, $f_{\pi N\Delta}$ and $g_{\rho\pi\pi}$ can be directly related to the experimental widths of $\Delta$ and $\rho$.

The $G_{\rho BB'}$ can also be readily related to the $f$’s and $g$’s defined in the Lagrangian formalism. Upon expressing the matrix elements $< B' | H_{\rho BB'} | B >$ of the Lagrangian model in a partial-wave representation and equating them to Eq. (3), we obtain at $p^2 = 0$

$$\langle Gv \rangle_{\rho NN} 4(\sqrt{2} - 1) \sqrt{3\pi^2 M_\rho M_N} = (fu)_{\rho NN} + (gu)_{\rho NN} = (1 + \kappa_\rho)(gu)_{\rho NN};$$

$$\langle Gv \rangle_{\rho N\Delta} \sqrt{\frac{3}{2}} (\sqrt{5} + 1) \sqrt{3\pi^2 M_\rho^3/M_\Delta} = (fu)_{\rho N\Delta},$$

where $Gv$ has the dimension of inverse mass, and $\kappa_\rho \equiv (f/g)_{\rho NN}$. The $u$ denotes the specific form factor employed in the Lagrangian model, which varies from one published work to
another. Using these two equations, one can make a straightforward comparison between the dipion and the MEP results. Our calculation gives the sum \((f + g)_{\rho NN}\). A knowledge of \(\kappa_\rho\) from the analysis of the nucleon electromagnetic form factor can be used to determine \(f\) and \(g\) separately.

For the calculation, the values \(g_{\rho\pi\pi} = 0.6684 M_\pi^{-1/2}\), \(\Lambda_{\rho\pi\pi} = 2.336 \text{ fm}^{-1}\), \(f_{\pi NN} = 1\), and \(\Lambda_{\pi NN} = \Lambda_{\pi N\Delta} = 950 \text{ MeV}(\equiv \Lambda(2))\) were used. The \(\rho\pi\pi\) parameters fit the \(\pi\pi\) phase shifts.\[9\] The value of the dipole range \(\Lambda(2)\) is taken from ref.[2]. It corresponds to a monopole range \(\Lambda(1)\).

\[\Lambda(2) = 950 \text{ MeV} \quad \text{(or \(\equiv \Lambda(2)\))} \]

For the other parameters, we used the following sets: (a) \(f_{\pi N\Delta} = 1.69\) (ref.[7]), \(f_{\pi\Delta\Delta} = 0.8\) (ref.[4]), \(\Lambda_{\pi\Delta\Delta} = \Lambda(2)\); (b) \(f_{\pi N\Delta} = 2.19\), \(f_{\pi\Delta\Delta} = 1.03\) (or \(f'_{\pi\Delta\Delta} = 1.55\)), \(\Lambda_{\pi\Delta\Delta} = \Lambda(2)\); and (c) \(f_{\pi N\Delta} = 2.19\), \(f_{\pi\Delta\Delta} = 1.03\), \(\Lambda_{\pi\Delta\Delta} = 1.86 \text{ GeV}\) corresponding to the extreme situation of having an equivalent monopole range \(\Lambda_{(1),\pi\Delta\Delta} = 1.2 \text{ GeV}\). The criteria leading to the above choices are as follows. In set (a) the ratio \(f_{\pi\Delta\Delta}/f_{\pi N\Delta} = \sqrt{2}/3\) satisfies a SU(6) quark model result.[11] Using the experimental \(\Gamma_{\Delta} = 115 \text{ MeV}\) to calculate \(f_{\pi N\Delta}\) and using \(g^2_{\pi\Delta\Delta}/4\pi = 60\) (given by the isobar analyses of the \(\pi N\rightarrow \pi\pi N\) data), we obtain set (b). As these isobar analyses do not give \(\Lambda_{\pi\Delta\Delta}\), we used \(\Lambda_{\pi\Delta\Delta} = 950 \text{ MeV}\) in set (b) and \(1.86 \text{ GeV}\) in set (c) in order to see the effects of this range parameter. We have also evaluated contributions by intermediate \(N^*(1440)\) state in the triangle diagrams, using \(f_{\pi N N^*} = 0.467\) and \(f_{\pi\Delta N^*} = 1.63\) calculated from the median values of experimental partial widths \(\Gamma_{N^*(\pi N)}\) and \(\Gamma_{N^*(\pi\Delta)}\). The full, coupled-equation results are presented in Table 1. The values within parentheses are due to the inclusion of the \(N^*\) contribution. The \((f + g)_{\rho NN}\) and \(f_{\rho N\Delta}\) have been evaluated, respectively, at \(w_c = 939\) and \(1232 \text{ MeV}\); they are found to be higher than the corresponding Born results by 5% and 14%. For the convenience of making comparisons, we have listed the values of \(f\) and \(g\), as obtained from Eqs.(6) and (7). Further, \(g_{\rho NN}\) has been calculated from the solutions \((f + g)_{\rho NN}\) with \(\kappa_\rho = f/g = 3.7\) given by an analysis of the electromagnetic form factor of the nucleon.

The MEP values of \(g_{\rho NN}\) vary from 1.28 to 2.27 (i.e. \(g^2/4\pi = 0.13 - 0.41\)). Those for \(f_{\rho N\Delta}\) vary from 4.91 to 7.81 (or \(f_{\rho N\Delta}^2/4\pi = 1.92 - 4.86\)). These lower and upper limits are, respectively, the results of refs.[3] and [7]. An inspection of Table 1 shows that with the inclusion of \(N^*\), sets (b) and (c) lead to a \(g_{\rho NN}\) that is very close to the MEP values. In
fact, a good agreement has been obtained when we used $f_{\pi NN}$, given by the experimental upper bound of $\Gamma_{N^*(\pi N)}$. While parameters of set (a) give a real part of $f_{\rho N\Delta}$ that is smaller than the lower limit of the MEP values by a factor of $\sim 2.5$, the use of $f_{\pi N\Delta}$ and $f_{\pi \Delta\Delta}$ derived directly from the data (set (b)) brings the difference down to within 40%. Use of a large $\Lambda_{\pi\Delta\Delta}$ (set (c)) makes the real part of $f_{\rho N\Delta}$ agree with the MEP values. However, in view of the deep inelastic scattering data on the upper limit of $\Lambda_{\pi NN}$ and $\Lambda_{\pi N\Delta}$, the large $\Lambda_{\pi\Delta\Delta}$ of set (c) may be questionable. Consequently, we consider the results due to set (b) as being more realistic. Our results also indicate that the contribution by $N^*$ is small. At this point, a comment on the complex nature of $f_{\rho N\Delta}$ is in order. As pointed out after Eq.(3), the coupling constants can be energy-dependent. An important feature of the di-pion model is that $f_{\pi N\Delta}$ becomes complex-valued at $w > M_{\pi} + M_{N} \equiv w_{th}$. The $f_{\rho N\Delta}$ of Table 1 were calculated at $w = M_{\Delta} > w_{th}$, where the $\pi N$ channel is open, thereby, giving rise to $\text{Im}[f_{\rho N;\Delta}] < 0$. The energy dependence of the real and imaginary parts of the $\rho N\Delta$ form factor at $p=0$ are shown, respectively, as the solid and dashed curves in fig.2, where we have defined $H(w, p) \equiv (Gv)(w, p)$.

Table 1: The coupling constants and ranges (in MeV) given by the di-pion model for the $\rho NN$ vertex at $w = M_{N}$ and for the $\rho N\Delta$ vertex at $w = M_{\Delta}$.

| Set | $(f + g)_{\rho NN}$ | $g_{\rho NN}$ | $\Lambda_{\rho NN}$ | $f_{\rho N\Delta}$ | $\Lambda_{\rho N\Delta}$ |
|-----|-----------------|-------------|-----------------|----------------|----------------|
| a   | 4.39(4.90)      | 0.934(1.04) | 435             | 1.89 - 2.09i(2.08 - 2.09i) | 495          |
| b   | 4.82(5.32)      | 1.025(1.13) | 430             | 3.50 - 2.71i(3.69 - 2.71i) | 490          |
| c   | 4.86(5.36)      | 1.033(1.14) | 430             | 5.25 - 2.70i(5.36 - 2.70i) | 490          |

Table 1 indicates that the ranges of the calculated form factors are in the region of 450 MeV, much smaller than the MEP values of 1.2 GeV. Here, we are considering the form factors as functions of $p^2$ and define $\Lambda$ as the momentum at which the magnitude of form factor is half of its value at $p^2 = 0$. Similar ranges have been obtained when we used the four-momentum transfer $p^2_{\rho}$ as the variable. We can, in fact, easily show that a smaller range is a consequence of the composite nature of a vertex.

We have further examined the four-pion content of the vertex functions by adding
the $\rho \to \omega \pi$ mechanism into the triangle diagrams. Using parameters of set (b) together with $g_{\omega \rho \pi} M_\omega = 9.71$ (ref. [13]), $g_{\omega N N} = 11.54$ (ref. [7]), and assuming the relation $g_{\omega \Delta \Delta} / g_{\omega N N} = g_{\pi \Delta \Delta} / g_{\pi N N}$ in our calculations, we have found that contributions from four-pion dynamics are negligible because $M_\omega \gg M_\pi$. Although diagrams having more than three subvertices could in principle also contribute, they tend to be unimportant because more are the particle propagators less is the interaction probability. It is, however, worth recalling that while both the $L = 1$ and $3$ interactions can contribute to the $\pi \Delta \Delta$ process, only the $L = 1$ case is considered in this and other works. Although inclusion of the $L = 3$ interaction can be easily implemented in the present formalism, experimental information on this f-wave coupling is sparse and uncertain. If one can establish experimentally that the f-wave coupling constant is not too small, then it could make a sizable contribution because the f-wave vertex function is $\propto q^3$ and $q$ is an integration variable that extends to very large values.

In summary, the di-pion model predicts $g_{\rho N N}$ that agrees with the MEP values, but the predicted $f_{\rho N \Delta}$ is about 30% smaller. In view of the fact that the $g_{\rho N N}$ and $f_{\rho N \Delta}$ of this work are based on a microscopic model while those of MEP were correlated fitting parameters, we regard the results given by the two approaches as being compatible. The inclusion of four-pion dynamics via the $\omega \pi$ doorway state does not alter the results, indicating that the di-pion mechanism does represent the leading mesonic contribution. However, experiments capable of providing information on f-wave $\pi \Delta \Delta$ couplings can help determine if mesonic dynamics alone is sufficient to account for all the strength of the phenomenological $\rho N \Delta$ coupling constants given by MEP. The ranges of the calculated form factors are in the 450 MeV region, much smaller than the corresponding MEP values, but in line with the small $\pi N N$ and $\pi N \Delta$ ranges. As mentioned earlier, a small range also reflects the composite nature of a vertex function. Smaller ranges will give calculated meson-exchange contributions that are different from those given by large ranges and, hence, may lead to new understandings of the data. Our analysis has shown that in the medium-energy regime, there is no compelling need for introducing non-hadronic dynamics. We stress, however, that MEP theory will become impractical at very high energies because the exchange of many heavy mesons has to be included in order to produce the correct energy dependence of the $N N$ cross section. Furthermore, the theory will be deficient because the quark substructure...
of the hadrons will start to manifest. We also stress that while it is a good approximation to employ real-valued vertex functions in analyzing $NN$ scattering below the pion production threshold,[7][10] the situation is different in meson production experiments where the energy $w_{th}$ can be surpassed. Thus, using real-valued vertex functions is questionable. We believe that a nonvanishing imaginary part of the form factor can give new interference effects in nuclear reaction calculations. This aspect of the meson-exchange dynamics merits a systematic investigation in the future.

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Figure 1:  Feynman diagrams for $\rho B B'$ vertex functions.
Figure 2: \(\text{Re}[H(w, 0)]\) (solid curve) and \(-\text{Im}[H(w, 0)]\) (dashed curve) of the Born \(\rho N\Delta\) vertex function versus \((w - M_N)/M_\pi\) obtained with parameters of set (b).
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