Peripheral nucleon-nucleon scattering is analysed in the framework of an effective field theory. Distorted-wave methods are used to remove the effects of one-pion exchange. Two-pion exchange and recoil corrections to one-pion exchange are then subtracted perturbatively. This removes all contributions up to order $Q^3$ in the chiral expansion. We have applied this to the $^1D_2$, $^1F_3$ and $^1G_4$ waves, using phase shifts from various partial-wave analyses by the Nijmegen group. In regions where these analyses agree we find no evidence for a breakdown of the chiral expansion. One of the terms in the effective short-range potential, the leading one in the $^1D_2$ channel, is larger than might be expected, but in general these terms have momentum scales of about 300–400 MeV. We also see hints of isospin breaking in the $\pi N$ couplings.

I. INTRODUCTION

Following the suggestion of Weinberg [1] and the pioneering work of Ordoñez, Ray and van Kolck [2], there has been much interest in using chiral perturbation theory (ChPT) to calculate long-range contributions to the nucleon-nucleon (NN) force. These have concentrated on two-pion exchange (TPE) [4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14], although work has also started on three-pion exchange [15, 16]. In an effective field theory like this, the residual short-distance interactions are represented by energy- or momentum-dependent contact interactions. These can be expanded in powers of the ratios between momenta or the pion mass and the scale of the physics which is responsible for these short-range interactions, for example the mass of the $\rho$ meson. Provided these ratios are small enough, this expansion should converge rapidly.

Peripheral nucleon-nucleon scattering in peripheral waves provides the best place to look for clear signals of the long-range TPE force. Various groups have done so [4, 5, 6, 10, 14] but they often find quite large deviations from the “experimental” phase shifts deduced from partial-wave analyses [17, 18]. This has led some to resort to introducing phenomenological cut-offs [6, 14] or additional degrees of freedom [5].

Recently the Nijmegen group has started carrying out PWA’s which include chiral TPE [7, 13] and they find that this makes significant improvements to their fits. The large differences between their results and available ChPT predictions have led to a claim that the fits cannot yield reliable values for parameters in the TPE potential [19]. However it is hard to make a direct comparison between the Nijmegen analyses and ChPT because they are done with a coordinate-space cut-off of 1.4 fm or larger. This means that short-distance parameters in the fits have to play two roles. As well as parametrising true short-range physics they also have to correct for artefacts introduced by the cut-off. Hence one cannot immediately tell from the Nijmegen parametrisation whether the residual short-range interactions are consistent with an effective field theory.

Distorted wave (DW) methods can be used to extract the effects of a known long-range interaction from two-body scattering, leaving a residual scattering amplitude which can be analysed using the techniques of effective field theory [20]. For systems with bound states close to threshold, such as $NN S$-waves, this is equivalent to a DW version of the effective-range expansion. In peripheral waves, where the scattering is weak, the effective short-range potential is essentially just an expansion of the residual $K$-matrix in powers of the energy.

Related approaches can be found in Refs. [4, 21] where a variable-phase method is used to construct the DW’s and a radial cut-off is imposed at some small radius. Ballot et al. [4] include TPE but do not attempt to extract a residual interaction strength, only examining the sensitivity of their results to the cut-off radius. (See also Ref. [10] for a similar treatment.) Pavon Valderrama and Ruiz Arriola [21] keep only OPE and parametrise the residual interaction via a boundary condition at their cut-off radius. This is constructed to reproduce exactly the effective-range expansion up to fourth order in the momentum.

Peripheral partial waves are of particular interest for testing $NN$ potentials from ChPT, since the nucleons do not

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1 For a recent review of effective theories of $NN$ scattering, see Ref. [3].
feel the very strong attraction that causes complications in the S-waves. However phase shifts in a wave with orbital angular momentum \( L \) typically grow like the \( L \)-th power of the energy. This means that quite small differences between the short-distance behaviour of potentials can lead to phase shifts that look very different at energies \( \sim 2m_r \). Conversely, at low energies, small differences between phase shifts can be hard to disentangle. In particular, differences between the available PWA’s in this region can be comparable in size to the TPE effects of interest, but this is not always obvious in simple plots of phase shifts. By applying the DW method to empirical phase shifts, we are able to obtain a residual interaction strength that is much less energy dependent than the phase shifts themselves. This has the advantage of providing a quantitative measure of the strength of the missing physics and of its energy dependence. It also makes much clearer the regions where the PWA results are not reliable.

Here we use the DW method to examine \( NN \) scattering in the \( ^1D_2, ^1F_3 \) and \( ^1G_4 \) partial waves using the phase shifts from five available Nijmegen analyses: PWA93, and the Nijmegen I, Nijmegen II, Reid93 and ESC96 potentials \[17\]. All of these include the long-range OPE potential but otherwise they parametrise the data quite differently. For example, the PWA imposes an energy-dependent boundary condition at a cut-off radius of 1.4 fm, while the potentials are energy-independent but may be either local or nonlocal. All of them have been fitted with similarly good \( \chi^2 \) to the \( NN \) data available in 1993, and so they can be regarded as alternative PWA’s. Any differences between these analyses should be taken as indications of the size of the systematic errors associated with them. The reason for concentrating initially on the spin-singlet waves is that the OPE potential is less singular than the centrifugal barrier and hence no regularisation is needed to construct the distorted waves. The need for some additional regulator in the triplet waves means that the short-range potential would have to play two roles, correcting for artefacts of the regulator as well as parametrising short-range physics, and so it would be harder to interpret.

The terms in the interaction can be classified according to a chiral expansion in powers of \( Q \), where \( Q \) denotes a factor of either a momentum or a pion mass \[3\]. In this counting the leading OPE potential is of order \( Q^0 \). We use the TPE potential given in Refs. \[4, 7\]. This includes terms up to order \( Q^3 \). As discussed by Friar \[8\], this potential should be used in conjunction with an OPE potential that has the usual nonrelativistic form (see Eq. \[2\] below) multiplied by \( M/E \) where \( E \) is the on-shell energy of one nucleon. This factor generates recoil corrections to OPE which start at order \( Q^2 \). Other corrections to OPE which might arise from higher-order \( \pi N \) vertices can all be absorbed in the on-shell \( \pi N \) couplings or \( NN \) contact terms (as discussed in Ref. \[3\]) or in higher-order terms in the expansion of \( M/E \).

The order-\( Q^2 \) recoil correction to OPE plus TPE terms from Ref. \[4, 7\] therefore provide the complete long-range potential at orders \( Q^2 \) and \( Q^3 \), with one exception: we use a single \( \pi N \) coupling constant. At this order there could be isospin breaking in the \( \pi N \) couplings \[22\], a point to which we shall return later. By using the DW method to extract all iterations of OPE and then subtracting these OPE and TPE terms, we are able to remove all contributions to the scattering up to order \( Q^3 \) in the chiral expansion.

## II. DISTORTED-WAVE METHOD

The starting point for the approach of Ref. \[20\] is the \( K \)-matrix which describes scattering between distorted waves of the known long-range potential, in our case OPE. On shell, this matrix \( K(p) \) is related to the observed phase shift \( \delta(p) \) by

\[
\hat{K}(p) = -\frac{4\pi}{M_p} \tan\left(\delta(p) - \delta_{\text{OPE}}(p)\right),
\]

where \( p \) is the on-shell relative momentum in the c.m. frame. Here \( \delta_{\text{OPE}}(p) \) denotes the phase shift for the lowest-order OPE potential. In the spin-singlet \( np \) channels this potential has the form

\[
V_{\text{OPE}}^{(0)}(r) = -f_{\pi NN}^2 \left[ e^{-m_\pi r} + 2 e^{-m_r r} \right],
\]

where \( m_\pi \) denotes the mass of the charged pion, \( m_\pi \) that of the neutral pion and the plus (minus) sign corresponds to isospin-singlet (-triplet) waves. We use the same value of the \( \pi N \) coupling as in the Nijmegen PWA’s \[17\], \( f_{\pi NN}^2 = 0.075 \). The waves are obtained by solving a Schrödinger equation with this potential. As expained in Ref. \[8\], relativistic effects can, to the order we are working, be absorbed into terms in the TPE potential and a factor of \( M/E \) multiplying the OPE potential. With standing-wave boundary conditions appropriate to the \( K \)-matrix, the DW’s have the asymptotic form

\[
\psi_{\text{OPE}}(p, r) \rightarrow \frac{\sin(pr - L\pi/2) + \tan\delta_{\text{OPE}}(p) \cos(pr - L\pi/2)}{pr}.
\]
Near the origin, they behave like
\[ \psi_{\text{OPE}}(r) \propto \frac{(pr)^L}{(2L+1)!!} \quad \text{as } r \to 0. \]

(4)

If the residual scattering is weak, we can represent it using an effective theory based on a trivial fixed point. The residual \( K \)-matrix \( \tilde{K}(p) \) is then equal to the leading matrix element of an energy-dependent short-range potential, in the DW Born approximation. A simple \( \delta \)-function potential will have no effect since the waves with \( L > 0 \) vanish at the origin. We could use an appropriate high derivative of a \( \delta \)-function to represent the short-range interactions but, for numerical implementation, it is more convenient to work with an energy-dependent \( \delta \)-shell potential,

\[ V_S(p,r) = \frac{[(2L+1)!!]^2}{4\pi R_0^{2L+2}} \tilde{V}(p) \delta(r-R_0), \]

(5)

Provided \( R_0 \) is chosen to be sufficient small that the asymptotic form (4) is valid, this form is numerically equivalent to a derivative of a \( \delta \)-function. (In practice we take \( R_0 = 0.1 \) fm.) Here we have divided out a factor of \( R_0^{-2L} \) so that the strength \( \tilde{V}(p) \) is independent of \( R_0 \) for small \( R_0 \). We have also divided out the numerical factor of \( [(2L+1)!!]^2 \) which is needed to compensate for the smallness of the high partial waves at small radii, as can be seen in Eq. (4).

The residual scattering when leading-order OPE only is removed starts at order \( Q^2 \) in the chiral expansion. Equating the DW matrix element of the short-range potential to \( \tilde{K}(p) \), we find that its strength is given by

\[ \tilde{V}^{(2)}(p) = \frac{R_0^{2L}}{[(2L+1)!! \psi_{\text{OPE}}(p,R_0)]^2} \hat{K}(p). \]

(6)

We can also remove the leading effects of order-\( Q^{2,3} \) OPE and TPE by subtracting from \( \hat{K}(p) \) the matrix elements of these potentials. The correction to OPE obtained by expanding the relativistic factor is

\[ V_{\text{OPE}}^{(2)}(r) = -\frac{p^2}{2M^2} V_{\text{OPE}}^{(0)}(r). \]

(7)

The forms of the TPE potentials can be found in Refs. [4, 7]. The resulting residual scattering starts at order \( Q^4 \) and can be described by a short-range potential with strength

\[ \tilde{V}^{(4)}(p) = \frac{R_0^{2L}}{[(2L+1)!! \psi_{\text{OPE}}(p,R_0)]^2} \left( \hat{K}(p) - \langle \psi_{\text{OPE}}(p) | V_{\text{OPE}}^{(2)} + V_{\text{TPE}}^{(2,3)} | \psi_{\text{OPE}}(p) \rangle \right). \]

(8)

The TPE potential at order \( Q^3 \) depends on the coefficients of three terms in order-\( Q^2 \) \( \pi N \) Lagrangian. In the results presented here, we have used the values obtained by the Nijmegen group from their recent analysis of \( N N \) data [13]: \( c_1 = -0.76 \, \text{GeV}^{-1} \), \( c_3 = -4.78 \, \text{GeV}^{-1} \), and \( c_4 = 3.96 \, \text{GeV}^{-1} \). We have also checked that other sets of values, for example those used in Refs. [4, 7, 15, 14], do not qualitatively change our results.

The leading short-distance interaction in a partial wave with orbital angular momentum \( L \) can be represented by a contact interaction proportional to the \( 2L \)-th derivative of a \( \delta \)-function. As mentioned above, we find it numerically more convenient to work with a \( \delta \)-shell interaction proportional to \( R_0^{-2L} \delta(r-R_0) \). In either case the leading interaction is of order \( Q^{2L} \) in the chiral expansion. In momentum space it has a strength of order \( \Lambda^{-2L+2} k^{2L} \) where \( \Lambda \) is a scale associated with the underlying short-distance physics. One might expect that for a potential of "natural" strength, the scale \( \Lambda \) should be of the order of the masses of the exchanged heavy mesons, at least 500 MeV. However, the rather strong \( \pi N \) coupling introduces another, significantly lower, scale of the order of 300 MeV. The current ChPT potentials do not include the \( \Lambda \), which also corresponds to a scale of about 300 MeV.

The leading OPE potential is of order \( Q^0 \) and so it has the form \( \Lambda^{-2} f_0(k/m_\pi) \) where \( k \) is a generic momentum variable and \( f_0(x) \) is a dimensionless function of order 1. TPE contributions start at order \( Q^2 \) and have the form \( \Lambda^{-4} m_\pi^2 f_2(k/m_\pi) \). If, as in the present work, we include OPE and TPE interactions up to order \( Q^3 \), the omitted order-\( Q^4 \) forces arising from exchange of up to three pions are of the form \( \Lambda^{-4} m_\pi^4 f_4(k/m_\pi) \).

The DW approach represents all higher-order effects in terms of short-range interactions. When we extract only the leading OPE potential, the effects of order-\( Q^2 \) TPE contributions in the partial wave \( L \) are replaced by an interaction with strength

\[ \tilde{V}^{(2)} \sim (L!)^2 \Lambda^{-4} m_\pi^{2(2L-1)} g_2(p/m_\pi), \]

(9)

where \( p \) is the on-shell momentum and \( g(x) \) is another dimensionless function of order 1. Here \( 2L \) powers of momenta have been extracted from \( f_2(k/m_\pi) \) to form a projector onto the relevant partial wave. This projector involves \( L \) derivatives of the initial and final wave functions, which leads to the numerical factor of \( (L!)^2 \).
When we extract OPE and TPE interactions up to order $Q^3$, the residual interaction strength has the form

$$\tilde{V}^{(4)} \sim (L!)^2 \Lambda^{-6} m_n^{-4-2L} k^{2L} g_4(p/m_n).$$

In the case of a $D$-wave, this contains a momentum-independent term $\Lambda^{-6}$, which corresponds to a contact interaction with an unknown coefficient.

By extracting the effects of the leading-order OPE and TPE forces, we expect to remove the dominant energy dependence of the scattering amplitudes up energies $T_{lab} \sim 200$ MeV. Beyond that region, three-pion exchange can start to contribute significantly, although calculations suggest that these forces are much smaller than other order-$Q^4$ contributions [15].

**III. RESULTS**

The main results of this analysis are shown in Figs. 1 and 2. (Note the differences in scale between the two panels of each plot.) There are a couple of general lessons to be drawn from them before examining the individual waves in more detail.

First, it is clear that for energies below about 80 MeV there are substantial differences between the various PWA’s. For the $^1F_3$ and $^1G_4$ waves in this region systematic artefacts of the different parametrisations completely dominate the differences between the empirical phase shifts and those from OPE plus TPE. Even in the range 150–250 MeV, the uncertainties are so large as to preclude much more than a rough estimate of the magnitude of the residual potential. Only in the $^1D_2$ wave does one find consistent results over a wide range of energies, from about 50 MeV upwards.
FIG. 3: The short-range effective potential in the $np\ ^1G_4$ partial wave, plotted in fm$^6$ against $T_{lab}$ in MeV. For other details see the caption to Fig. 1.

FIG. 4: The short-range effective potential (in fm$^6$) in the $np\ ^1G_4$ partial wave, with leading-order OPE and order-$Q^2$ TPE removed but not the order $Q^3$ recoil correction to OPE. For other details see the caption to Fig. 1.

One point to note about the $^1D_2$ and $^1G_4$ waves is that we find no correlation among the deviations of the different PWA’s from their common trend. This suggests that there is no systematic bias to the fits.

Second, it is important to use the “correct” $\pi N$ coupling, namely the one assumed in the PWA, and to include the $M/E$ factor multiplying OPE (or at least the leading correction from it). If these are not done then the results show a strong systematic energy dependence at low energies. An example of this is given in Fig. 3 where TPE has been subtracted but not the order-$Q^2$ correction to OPE. There are plots in Refs. [4, 6, 10] comparing OPE plus TPE with PWA’s that all look similar to each other. However Ref. [4] includes the $M/E$ factor but uses a large value for the $\pi N$ coupling ($f_{NN} = 0.077$ as opposed to the Nijmegen recommended value). In contrast Refs. [6, 10] omit the $M/E$ factor but use smaller $\pi N$ couplings. The results of these two choices are quite similar for $^1F_3$ and $^1G_4$ scattering in the energy range 100–200 MeV, where difference from pure OPE are most visible. In both cases, the differences from the PWA’s are larger than those obtained when the consistent $\pi N$ coupling is used and recoil corrections are included. Although these differences are similar in magnitude to the systematic uncertainties in the current PWA’s, they will become more significant when phase shifts from improved PWA’s become available [5, 13].

In the $^1D_2$ wave, shown in Fig. 1 subtraction of the order-$Q^2$ terms may not have much effect on the size of the residual scattering, but it does dramatically reduce its energy dependence. There is a nearly 100% change over the energy range 50 to 300 MeV if OPE alone is removed, but this is reduced to about 20% when the order-$Q^2$ terms are subtracted. The typical size of the residual strength after removal of OPE, $V^{(2)}$, is about 0.1 fm$^6$. Comparing this with Eq. (9) we find that the corresponding scale $\Lambda$ is approximately 300 MeV, as expected for pion-exchange forces. After subtracting OPE and TPE to order $Q^3$, the residual short-range potential is, to a very good approximation, linearly dependent on the energy. Its intercept occurs at about $\tilde{V}^{(4)} \approx 0.16$ fm$^6$, which corresponds to a momentum scale of about 200 MeV in Eq. (10). This is distinctly smaller than one would expect on grounds of “naturalness”. Its slope is $d\tilde{V}^{(4)}/dp^2 \approx 0.012$ fm$^8$, corresponding to a scale of about 370 MeV.

This picture bears out what was found in Refs. [4, 10], namely that one counterterm of order $Q^4$ is able to explain the bulk of the residual scattering after all contributions up to order $Q^3$ have been removed. One might worry that
those authors neglected the recoil correction to OPE. However in this channel the TPE contributions are so much larger that the neglect of recoil does not affect the results significantly. At least for the spin-singlet channels, this analysis also removes the worries raised in Ref. [14]: although there is one unnaturally large term in the effective short-range potential, there is no evidence for a breakdown of the chiral expansion, even up to energies of \( \sim 300 \text{ MeV} \).

In the \( ^1F_3 \) and \( ^1G_4 \) cases there are large differences between the various PWA’s for energies below about 150 MeV and so it is hard to draw such definite conclusions. Nonetheless it is clear that the residual scattering is much smaller after the order-\( Q^{2,3} \) terms have been subtracted, as in the results of Refs. [14, 15]. The sizes of the residual scattering strengths, \( V^{(2)} \) and \( V^{(4)} \), at energies below 100 MeV in these waves correspond to scales in the region 300–400 MeV.

Although the different PWA’s show no overall bias in the \( ^1D_2 \) and \( ^1G_4 \) waves, this is not the case in the \( ^1F_3 \) wave. There, in contrast, all the residual interactions show a significant downward curvature at low energies, even after subtraction of the order-\( Q^{2,3} \) terms. While this may just reflect the fact this wave is poorly constrained by data (there are differences between the Nijmegen [17] and VPI [18] phase shifts for this wave) a more intriguing possibility is that it could be a signal of a long-ranged isospin-breaking effect.

In this context, it should be noted that the isospin-triplet partial waves, such as \( ^1D_2 \) and \( ^1G_4 \), are fitted to \( pp \) scattering data in the Nijmegen PWA’s [17]. The corresponding \( np \) results are then obtained by simply replacing the neutral-pion exchange in \( pp \) by the relevant combination of neutral- and charged-pion exchange, as in Eq. (2).

In contrast the isospin-singlet waves, such as \( ^1F_3 \), have to be obtained from fitting \( np \) data. A single value for the \( \pi N \) coupling is used in all cases. The overall downward deviation in our results is consistent with what would be expected if we had taken too high a \( \pi N \) coupling in the DW analysis of these waves. A similar pattern is also seen in the \( ^1P_1 \) wave. The size of the effect is compatible with what can be deduced from PWA’s of \( NN \), \( \pi N \) or \( N\bar{N} \) scattering [24], however we would caution that all these effects are comparable to the uncertainties in the PWA’s. Other isospin-violating \( NN \) interactions in the framework of ChPT have been discussed in Refs. [16, 25, 26]. The longest-ranged of these, and hence the most important at low-energies, are the electromagnetic corrections to OPE [25], but we find that the effect of subtracting the order-\( \alpha \) corrections to OPE is very small for energies above 80 MeV. It will be interesting to see if this deviation remains when the DW method is applied to phase shifts from the newer chiral PWA’s of the Nijmegen group [7, 13].

**IV. CONCLUSIONS**

We have presented here a method for extracting the effects of OPE and TPE from peripheral \( NN \) phase shifts. A DW approach is used to remove the effects of simple OPE to all orders. TPE and recoil corrections to OPE are then subtracted perturbatively and we are then able to remove all contributions up to order \( Q^3 \). We have applied this technique to peripheral \( NN \) scattering in spin-singlet waves, using phase shifts from various Nijmegen PWA’s [17]. We find residual interactions which are consistent with the expectations for an effective field theory. In the \( ^1D_2 \) wave, the energy dependence of the residual interaction is essentially linear up to nearly 300 MeV. The systematic errors in the various PWA’s are large in the \( ^1F_3 \) and \( ^1G_4 \) waves, making it hard to draw definite conclusions about the effectiveness of the theory in these cases. Nonetheless, the residual scattering in these waves is small after removal of all order-\( Q^{2,3} \) terms.

The momentum scales of the terms in the residual short-range potentials are, with one exception, 300 MeV or larger, as expected from the scale appearing the pion-exchange potentials. The exception is the leading term in the \( ^1D_2 \) wave. This term is unnaturally large, corresponding to a scale of about 200 MeV. In the energy region up to \( T_{\text{lab}} \sim 200 \text{ MeV} \), where OPE and TPE forces are expected to dominate, there is no evidence of a breakdown of the chiral expansion, nor is it necessary to introduce any additional regularisation in the spin-singlet channels.

In the isospin-singlet waves we find hints of isospin breaking in the \( \pi N \) couplings, although these are very much at the limit what can be deduced given the systematic errors in the different PWA’s.

It will be very interesting to see the results of this DW method when it is applied to phase shifts from the newer Nijmegen PWA’s [7, 13]. If these are sufficiently well-determined at low-energies, it should be possible to use the method to examine whether the results after extraction of the order-\( Q^3 \) potential [11, 12, 13] remain consistent with the chiral expansion.
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