Parity nonconserving proton-proton elastic scattering

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Abstract

The parity nonconserving longitudinal analyzing power $A_L$ is calculated in elastic $\bar{p}p$ scattering at the energies below the approximate inelastic region $T_{\text{lab}} = 350$ MeV. The short-ranged heavy meson $\rho$ and $\omega$ exchanges as well as the longer-ranged $2\pi$ exchanges are considered as the mediators of the parity nonconserving interactions. The DDH ”best” coupling values are used as the parity nonconserving meson-NN couplings. Also three different parity nonconserving two-pion exchange potentials by various authors are compared.

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I. INTRODUCTION

Weak interaction is distinct in the leptonic, semileptonic, and strangeness nonconserving hadronic processes. However, it is not so clear-cut in the strangeness conserving hadronic sector due to its diminutive strength against that of incessantly present strong interaction. Nevertheless, the parity nonconserving (PNC) weak interaction is unique in the sense that it sorts out different helicity states unlike any other interaction. For this particular reason, it can, in principle, be extracted under those overwhelming and unfailingly parity conserving (PC) strong and electromagnetic interactions.

Even though a direct heavy $Z^0$ or $W^\pm$ boson exchange is highly improbable over the internuclear distances, it is feasible between the nucleon and virtual meson. Consequently, the PNC $NN$ interactions may be parametrized by weak meson-$NN$ coupling constants modelled in terms of quarks and intermediate bosons. Traditionally the PNC $NN$ calculations have relied largely on the single meson exchange picture, based on the DDH potential [1] in which the PNC $NN$ interactions are due to $\pi^\pm$, $\rho$, and $\omega$ exchanges. Nowadays at very low energies, the calculations are preferably done in the framework of the model-independent effective field theories (EFT). However, all these models are parameterized by about half a dozen weak meson-$NN$ couplings (see e.g. Refs. [2, 3]), which are, even today, insufficiently known despite all the experimental and theoretical efforts.

Due to the fact the PNC interactions treat unequally different helicity states, the PNC nucleon-nucleon ($NN$) experiments are inherently based on the spin control of the particle. Probably the cleanest observable, in the sense that it is nearly a 100% pion exchange dominated, arises from the radiative PNC reaction $\bar{n}p \rightarrow \gamma d$ at threshold. The ongoing NPDGamma experiment [4] aims to determine the weak $\pi NN$-coupling $h_\pi^{(1)}$ by measuring the $\gamma$-asymmetry of this reaction, with such an accuracy that should elucidate the correctness of the most preferred value of the $h_\pi^{(1)} = 4.6 \times 10^{-7}$ suggested by DDH. Instead, the PNC $\bar{\gamma}d \leftrightarrow np$ reactions at threshold would lead only to nonpionic exchange effects, despite of small exchange currents and $\Delta$-effects [5]. In any case, when it comes to elastic PNC $\bar{p}p$ scattering, it is generally believed (based on the simple single meson exchange picture) that the pion does not contribute to it due to the lack of $\pi^0$-exchange. This is because in general the PNC neutral spinless meson exchange, e.g. $\pi^0$, is forbidden by the simultaneous violations of the P and CP symmetries [6]. However, not only the fact that the strong and
weak (DDH) pion couplings are sizable, but also that the pions are nearly six times lighter than heavy mesons, it seems reasonable to assume that the longest-range and possibly the leading effects are nonetheless due to pion exchange (in this particular case, induced by the two charged pions).

For the measurement of the PNC $\vec{p}p$ longitudinal analyzing power $\bar{A}_L$, there exist three precision experimental data points: Bonn at 13.6 MeV $(-0.93 \pm 0.21) \times 10^{-7}$, PSI at 45 MeV $(-1.50 \pm 0.22) \times 10^{-7}$, and TRIUMF at 221.3 MeV $(0.84 \pm 0.29) \times 10^{-7}$. The Bonn and PSI experiments are low energy scattering experiments, where the contribution to the $\bar{A}_L$ arises only from the lowest $^1S_0 - ^3P_0$ transition. The TRIUMF experiment, on the contrary, is a transmission experiment with the energy chosen so that the contribution arises merely from the $^3P_2 - ^1D_2$ and $^1D_2 - ^3F_2$ transitions. At the energy of the TRIUMF experiment, the $^1S_0$ and $^3P_0$ phases serendipitously cancel out due to strong interaction interference from which follows that the $J = 0$ transition goes to zero, while the $J = 4$ and higher ones still remain insignificant. What is more, for $J = 2$, the local and nonlocal contributions of the $\omega$ exchange mostly cancel out because of a small isoscalar anomalous magnetic moment $\chi_\omega$. In contrast for the $\rho$ exchange, the local contributions dominate over the nonlocal ones because of a large isovector anomalous magnetic moment $\chi_\rho$.

Assuming that the $J = 2$ mixing arises from the $\rho$ exchange, the central goal of the TRIUMF experiment was to determine the weak $\rho pp$-coupling $h_{\rho pp} = h_{\rho}^{(0)} + h_{\rho}^{(1)} + h_{\rho}^{(2)}/\sqrt{6}$ whereas the lower-energy experiments Bonn and PSI determined the $h_{\omega}^{pp} + h_{\omega}^{pp}$, where $h_{\omega}^{pp} = h_{\omega}^{(0)} + h_{\omega}^{(1)}$. In these experiments the reasoning was built on the DDH potential. However, already the work [10] including the effect of intermediate $N\Delta(1232)$ states via $\rho$ exchanges in the coupled channels showed that the simplest and most straightforward interpretation of the TRIUMF experiment might not be enough. The $\Delta$ effect was significant enough to suggest that the coupling could rather be effective involving $\rho$ exchange both in $NN$ and $NN \leftrightarrow N\Delta$ transitions. In our later work Ref. [11] on PNC $\vec{p}p$ elastic scattering, we looked at the effects of the $N\Delta$-channels in the coupled-channels formalism as well as the effects of the two-pion exchange (TPE). The effects were again found significant and cast doubt on the aforementioned $h_{\rho pp}$-coupling and whether its value is straightforwardly proportional to the TRIUMF data point. The preceding works Refs. [12, 13] on the reaction in question take into account the TPE, of which the former investigates it as a part of the short-ranged $\rho$ meson exchange and the latter considers it in the framework of the EFT.
As for the present work, we should stress that the purpose of this paper is no more
than to emphasize the importance of the TPE in the PNC $\bar{p}p$ elastic scattering, which
should be clear from the related model dependencies. Even though the TPE is far more
complicated than the single meson exchange, it should not be ignored in this particular
case due to its considerable strength and range. As shown in Ref. [11], another possibly
noteworthy contribution arises from the $\Delta$-resonance even at low energies, but it is not taken
into consideration here in its fullest form because of the large uncertainties related to the
meson-$N\Delta$ couplings especially in the weak sector. The $\Delta$ is taken into account only to the
extent it appears in the PNC TPE potentials. Since there is no $\pi N\Delta$ coupling related to a PNC vertex [14] or it is small [15, 16] (we take it as zero), then on the side of the weak couplings, the PNC TPE effects are only proportional to the $h^{(1)}_{\pi}$. Besides the DDH, there are various calculations [14, 17–23] for the $h^{(1)}_{\pi}$ (ranging between 0 and $3.4 \times 10^{-7}$) indicating a smaller value than what is the DDH ”best” recommendation. The hope is that the NPDGamma experiment would reduce the obscurity of this coupling constant.

This work is based on the use of the distorted-wave Born approximation (DWBA) and
the optical theorem. In the calculations, we employ the Reid93 potential [24] taking into
account the lowest five parity admixed transitions, i.e. the total angular momentum up to
$J = 4$. The short-ranged contributions are taken as the results of heavy meson $\rho$- and $\omega$-
exchanges, for which we use the DDH potential. For the long-ranged effects, we compare
three different PNC TPE potentials on the market given in Refs. [11, 25, 26]. Note that,
besides Ref. [26], there exists also another chiral perturbation theory (ChPT) derivation for
the PNC TPE $N\Delta$ potential [27], which however is not utilized in this work.

The remainder of the paper is outlined as follows. Section II gives the basic formalism
for the calculation of the PNC $\bar{p}p$ elastic scattering and Sec. III summarizes the results.

II. FORMALISM

The PNC $\bar{p}p$ elastic scattering experiments measure the difference between the cross-
sections $\sigma_{m_1}$ of the transmitted protons with the spins parallel ($m_1 = \frac{1}{2}$) and antiparallel
($m_1 = -\frac{1}{2}$) along the direction of propagation. The PNC analyzing power is given as

$$A_L = \frac{\sigma_{\frac{1}{2}} - \sigma_{-\frac{1}{2}}}{\sigma_{\frac{1}{2}} + \sigma_{-\frac{1}{2}}} \quad (1)$$
where in the other words $\sigma_{1/2}$ and $\sigma_{-1/2}$ denotes respectively the total cross-sections of the positively and negatively polarized proton beam.

The complete potential for the PNC $pp$ interaction is a combination of distortive PC potentials and small perturbatively treated PNC potentials $\hat{V} = \hat{V}^{PC} + \hat{V}^{PNC}$. The PC potential is the sum of the Coulomb $\hat{V}^{PC}_C$ and nuclear $\hat{V}^{PC}_N$ $NN$ potentials, where we take $\hat{V}^{PC}_N$ as the Reid93 potential [24]. The PNC potential is considered to arise from the long-ranged TPE potential and short-ranged heavy meson potential $\hat{V}^{PNC} = \hat{V}^{PNC}_2 + \hat{V}^{PNC}_{\rho,\omega}$. The used PNC TPE potentials [11, 25, 26] are respectively abbreviated by the authors as NPI, DHAL, and K. The potentials $\hat{V}^{DHAL}_{2\pi}$ and $\hat{V}^{K}_{2\pi}$ are built on QCD based ChPT and the $\hat{V}^{NPI}_{2\pi}$ on the time-ordered perturbation theory. The $\hat{V}^{DHAL}_{2\pi}$ results essentially from the $v^{EFT}_{44}(q)$ in Eq. 12 of Ref. [25]. Notable is that it comprises only the $NN$ intermediate states while $\hat{V}^{K}_{2\pi}$ and $\hat{V}^{NPI}_{2\pi}$ include also the $N\Delta$ intermediate states. Anyhow, they all are the spin changing local potentials of the form $\hat{V}^{PNC}_{2\pi}(r) = h_{\pi}^{(1)}(\hat{\tau}_1z + \hat{\tau}_2z)(\sigma_1 \times \sigma_2) \cdot rW(r)$ in the two-proton case. The radial functions $W(r)$ have different structures in each potential. The $NN$ parts of the unregularized $\hat{V}^{K}_{2\pi}(r)$ and $\hat{V}^{DHAL}_{2\pi}(r)$ potentials are identical, apart from the $\delta(r)$-function term in the latter one arising from the constant term in momentum space in its dispersion relation. However, since we are not only dealing with low energies, these two potentials should be provided with form factors, in which case they differ from each other even if regularized by the same form factors. The DDH "best" value $h_{\pi}^{(1)} = 4.6 \times 10^{-7}$ is used in the PNC TPE potentials.

As a PNC heavy meson potential, we use the DDH potential and their "best" weak meson-$NN$ coupling values. The isospin matrix element of the DDH potential, taken between the initial and final $pp$ states, is

$$\hat{V}^{PNC}_{\rho,\omega}(r) = - \sum_{\alpha=\rho,\omega} g_\alpha h^{pp}_\alpha \frac{1}{2M} \left( (\sigma_1 - \sigma_2) \cdot \{-i\nabla, Y_\alpha(r)\} + i(1 + \chi_\alpha)(\sigma_1 \times \sigma_2) \cdot [-i\nabla,Y_\alpha(r)] \right),$$

(2)

with $h^{pp}_\rho = h^{(0)}_\rho + h^{(1)}_\rho + h^{(2)}_\rho / \sqrt{6}$ and $h^{pp}_\omega = h^{(0)}_\omega + h^{(1)}_\omega$, which have the numerical values of $-15.48$ and $-3.00$ in units of $10^{-7}$ respectively. As for the other parameters, we take the values for the strong couplings as $g_\rho = 2.79$ and $g_\omega = 8.37$ and for the anomalies as $\chi_\rho = 3.71$ and $\chi_\omega = -0.12$. The radial $Y_\alpha(r) = \exp(-m_\alpha r)/4\pi r$ are the Yukawa functions, which we use only in the form modified by the dipole form factors of the type $(\Lambda^2_\alpha - m^2_\alpha)^2(q^2 + \Lambda^2_\alpha)^{-2}$ taking the cut-off masses as $\Lambda_\rho = 1.3$ GeV and $\Lambda_\omega = 1.5$ GeV. The $\hat{V}^{NPI}_{2\pi}$ potential is slightly
scaled to correspond to the strong $\pi NN$ coupling value of $g_\pi = 13.45$, which is in line with the other couplings. The heavy meson masses are $m_\rho = 770$ MeV and $m_\omega = 782$ MeV, and $M = 939$ MeV is the average nucleon mass. We call the couplings given above as the standard set, since they are the typical choice in the PNC calculations.

The $pp$ scattering amplitude $f(k, \theta) = f^C(k, \theta) + f^N(k, \theta)$ consists of the Coulomb scattering amplitude (superscripted by C) representing electromagnetic interaction and the Coulomb-nuclear scattering amplitude (superscripted by N) including electromagnetic, strong, and weak interactions. The Coulomb scattering amplitude is given by

$$f^C(k, \theta) = -\frac{\eta}{2k \sin^2 \frac{\theta}{2}} e^{i[2\sigma_0 - \eta \ln \sin^2 \frac{\theta}{2}]}$$

where $\eta = \alpha \mu / k$, $\alpha$ is the fine-structure constant, $\mu = M/2$ is the reduced mass of the two nucleons, and $\sigma_0 = \arg \Gamma(1 + i\eta)$ is the Coulomb S-wave phase shift. An awkward feature of Eq. (3) is that it is undefined at $\theta = 0$. Thus, in the determination of the total scattering cross-section by means of the optical theorem, the singularity of the total scattering amplitude in the forward direction is simply removed by the subtraction of the $f^C(k, 0)$, leaving only the $f^N(k, 0)$ to contribute. The forward, $\theta = 0$, $\bar{pp}$ Coulomb-nuclear scattering amplitude in the DWBA is given by

$$f^N_{m_1m_2}(k, 0) = -\frac{\mu}{2\pi} \left[ C\langle k \hat{z}; m_1m_2 | V^{PC}_N | k \hat{z}; m_1m_2 \rangle^{(+)} + (-) \langle k \hat{z}; m_1m_2 | V^{PNC}_N | k \hat{z}; m_1m_2 \rangle^{(+)} \right],$$

where the nuclear potentials are sandwiched between the Coulomb-distorted strong interaction wavefunctions. The $pp$ wavefunctions are of the form

$$\langle r | k \hat{z}; m_1m_2 \rangle^{(\pm)} = \sum_{SM_S} \frac{1}{\sqrt{8\pi}} \sum_{LJ} i^L \sqrt{2L + 1} \langle L0SM_S | JM_S \rangle e^{\pm i\sigma_L} U_{LL'}^{SJ}(k, r) \mathcal{Y}_{JM_S}^{L}\hat{r} | 11 \rangle,$$

where the z-axis is taken along the direction of $k$, $\mathcal{Y}_{JM_S}^{L}\hat{r}$ are the eigenfunctions of the coupled angular momentum, and $|11\rangle$ denotes the isospin state $|TM_T\rangle$ of the two-protons. In the wavefunctions with the subscript C, as it is in final state of the PC amplitude, the radial wavefunctions $U_{LL'}^{SJ}(k, r)$ (including the phase shifts) are simply replaced by the
regular Coulomb functions $F_L(kr)$, which further reduce to the spherical Bessel functions $j_L(kr)$ if the Coulomb interaction is turned off, i.e. $\eta = 0$. From the given equations, the longitudinal scattering asymmetry becomes

$$A_L(k) = \frac{\text{Im} \sum_{SS'} (-) \langle k \hat{z}; S'0|\hat{V}_{\text{PNC}}| k \hat{z}; S0 \rangle^{(+)} }{\text{Im} \sum_{SM_S} c \langle k \hat{z}; SM_S|\hat{V}_{\text{PC}}| k \hat{z}; SM_S \rangle^{(+)} }.$$  \hfill (7)

While the lower energy experiments measure directly the scattered particles, the TRIUMF E497 and higher energy experiments measure the transmitted beam after passing through the target, see e.g. Ref. [28] for a summary of the PNC $\vec{p}\vec{p}$ experiments. In transmission experiments, a complication arises due to the fact the Coulomb interaction is singular in the forward direction. Therefore, we consider the Coulomb distortions near the propagation direction of the transmitted beam, as done, e.g. in Refs. [29] and [30]. Symmetrized and properly normalized Coulomb scattering amplitude may be written as

$$f_{m_1m_2}^C(k, \theta) = \frac{1}{\sqrt{2}} \sum_{SM_S} \langle \frac{1}{2} m_1 \frac{1}{2} m_2 | SM_S \rangle \langle \frac{1}{2} m'_1 \frac{1}{2} m'_2 | SM_S \rangle \left[ f^C(k, \theta) + (-)^S f^C(k, \pi - \theta) \right],$$  \hfill (8)

where $f^C(k, \theta)$ is given in Eq. (3). The spin averaged Coulomb cross-section for a transmission experiment becomes

$$\sigma_{m_1}^{C\theta_0}(k) = \pi \sum_{m_1m_2} \int_{\theta_0}^{\frac{\pi}{2}} d\theta \sin \theta |f_{m_1m_2}^C(k, \theta)|^2 = \frac{\pi \eta^2}{2k^2} \left( \frac{1}{\sin^2 \frac{\theta_0}{2}} - \frac{1}{\cos^2 \frac{\theta_0}{2}} + \frac{1}{\eta} \sin[2\eta \ln \tan \frac{\theta_0}{2}] \right) \delta_{m_1m_1},$$  \hfill (9)

where $\theta_0 > 0$ is such a small cut-off angle that $f^N(k, \theta_0) \approx f^N(k, 0)$. The corresponding nuclear cross-section is

$$\sigma_{m_1}^{N\theta_0}(k) = \sigma_{m_1}^{N}(k) - 2\pi \int_0^{\theta_0} d\theta \sin \theta \frac{d\sigma_{m_1}^{N}}{d\Omega}(k, \theta) = \frac{\pi}{k} \sum_{m_2} \text{Im} \left( f_{m_1m_2}^N(k, 0)e^{2i[\eta \ln \sin \frac{\theta_0}{2} - \sigma_0]} \right),$$  \hfill (10)

where $\sigma_{m_1}^{N}(k)$ is the total cross-section given by the optical theorem and the differential cross-section is taken as $d\sigma_{m_1}^{N} = d\sigma_{m_1} - d\sigma_{m_1}^{C}$. In the last step of Eq. (10), the result

$$\int_0^{\theta_0} d\theta \sin \theta f^C(k, \theta) = \frac{1}{ik} \left( 1 - e^{2i[\eta \ln \sin \frac{\theta_0}{2} - \sigma_0]} \right),$$  \hfill (11)

first derived in Ref. [31], was used. The longitudinal transmission asymmetry becomes

$$A_{L0}^\theta(k) = \frac{\text{Im} \left[ \sum_{SS'} (-) \langle k \hat{z}; S'0|\hat{V}_{\text{PNC}}| k \hat{z}; S0 \rangle^{(+)} e^{2i[\eta \ln \sin \frac{\theta_0}{2} - \sigma_0]} \right] }{\text{Im} \left[ \sum_{SM_S} c \langle k \hat{z}; SM_S|\hat{V}_{\text{PC}}| k \hat{z}; SM_S \rangle^{(+)} e^{2i[\eta \ln \sin \frac{\theta_0}{2} - \sigma_0]} \right] - \frac{4k}{M} \sigma_{m_1}^{C\theta_0}(k) }.$$  \hfill (12)
III. RESULTS

Now we give the results at energies ranging between 1 and 350 MeV for the PNC longitudinal analyzing powers. In all cases the Reid93 potential is employed and the shown experimental data points are the Bonn, PSI, and TRIUMF ones, for which the values are given in the introduction. In Figs. 1-3 we use the standard set of couplings given in Sec. II whereas in Fig. 4 we use a weaker πNN coupling together with the heavy meson couplings of the configuration space Bonn potential [32]. In all Figs. 1-4 we employ the \( \hat{V}_{2\pi}^{\text{NPI}} \) as a reference TPE potential along with the PNC heavy meson exchange potential Eq. (2) with the dipole form factors of the type \((\Lambda_{\alpha}^2 - m_{\alpha}^2)^2(q^2 + \Lambda_{\alpha}^2)^{-2}\) with the cut-offs of \(\Lambda_{\rho} = 1.3\) GeV and \(\Lambda_{\omega} = 1.5\) GeV. Basically, the PNC pp effects are exclusive properties of nuclear interactions disturbed by the Coulomb field within the range of the nuclear forces. To obtain a clean PNC signal, the external long-range Coulomb effects can be cut out. In all figures, we utilize the scattering analyzing power of Eq. (7), in which the long-range Coulomb effects are neglected by omitting the Coulomb phases \(e^{i\sigma L}\) of the wavefunctions in Eq. (6). However, these negligible effects are included in the asymmetries of Fig. 3 where also the transmission analyzing power of Eq. (12) is depicted.

Figure 1 shows separately the \(\rho-, \omega-, \) and TPE contributions to the asymmetry. Throughout the energy range, the TPE effect is about twice as large as that of the heavy mesons and, as a consequence, the calculated asymmetry sets within the error limits of the experimental data. Figure 2 depicts the contributions of the different parity admixed partial waves

![Figure 1. The contributions of the \(\rho-, \omega-, \) and \(2\pi-\) exchanges to the analyzing power.](image)
up to $J = 4$. The transitions with $J = 4$ (or higher) are unimportant and, thus, the lowest three admixtures would in fact be sufficient within the used energy range. One particularly interesting feature of the asymmetry, as was first pointed out in Ref. [33] and utilized in the TRIUMF experiment, is that the $^1S_0 - ^3P_0$ contribution vanishes at a specific energy due to the equal, but opposite phase shifts of the $^1S_0$ and $^3P_0$ partial waves, which is seen at 224.7 MeV in Fig. 2.

The long-range Coulomb effects to the asymmetries are illustrated in Fig. 3 along with the cut-off angle $\theta_c$ dependence of the transmission asymmetry. The calculated scattering
and transmission asymmetries at the energies of about 150 MeV and above become nearly indistinguishable by the angles $\theta_c \geq 2^\circ$. Especially noteworthy is that at the energy of the only transmission experiment, TRIUMF, the asymmetry remains practically unaffected. Just to show the strong coupling sensitivity to the analyzing power, in Fig. 4 we employ an alternative set of couplings, $g_\pi^2/4\pi = 13.8$ for the pion and the Bonn potential configuration space values $g_\rho^2/4\pi = 0.95$, $g_\omega^2/4\pi = 20$, $\chi_\rho = 6.1$, and $\chi_\omega = 0$ for the heavy mesons. Compared to the use of the standard set of couplings, the asymmetry is enhanced by this choice of couplings. Also the TPE and heavy meson exchange contributions to the analyzing power become about equal. In the same figure, we have separated the nonlocal and local contributions of the PNC heavy meson exchange potential, which arise respectively from the anticommutator and commutator terms of Eq. (2). The resulting curves are formally consistent with Ref. [29]. Note that the scaling between Figs. 1 and 4 is straightforward for $\rho + \omega$ total and TPE, since only the aforementioned strong couplings are changed.

Figures 5 and 6 represent respectively the PNC TPE potential $\hat{V}_{2\pi}^{K}$ and $\hat{V}_{2\pi}^{DHAL}$ contributions to the scattering asymmetry. Because the PNC potentials are in general treated perturbatively in the DWBA, the regularization of them is not vital. However, since the effect of the singularity comes forth more and more along with the increasing energy, it should be removed by the regularization as usual. In contrast, the chiral perturbation the-
FIG. 5. The TPE by the $\hat{V}_{2\pi}^K(r)$ potential with the monopole (F) and dipole (FF) form factors using the cut-off masses $\Lambda = 1.0$ GeV and 1.2 GeV and also without (w/o) the form factors. As a comparison, the asymmetry is also plotted using the $\hat{V}_{2\pi}^{NPI}(r)$ potential in which the coupling values are scaled to correspond to those of the $\hat{V}_{2\pi}^K(r)$.

ory based $\hat{V}_{2\pi}^K$ and $\hat{V}_{2\pi}^{DHAL}$ potentials would serve their purpose best as unregularized due to their model independent nature. However, in Figs. 5 and 6 these potentials are used both with (F or FF) and without (w/o) regularization. When regularized, we incorporate the monopole $\Lambda^2(q^2 + \Lambda^2)^{-1}$ (F) and dipole $\Lambda^4(q^2 + \Lambda^2)^{-2}$ (FF) form factors using two different cut-off masses $\Lambda = 1.0$ GeV and $\Lambda = 1.2$ GeV. A monopole form factor of the same type is also used to the $np$ part of the $\hat{V}_{2\pi}^{DHAL}$ potential in the radiative reaction $\vec{n}p \rightarrow \gamma d$ in Ref. [34]. As seen in Figs. 5 and 6 the resulting asymmetries are in most cases formally similar for all the $\hat{V}_{2\pi}^K$, $\hat{V}_{2\pi}^{DHAL}$, and $\hat{V}_{2\pi}^{NPI}$ potentials. When switching over from the monopole to dipole type form factor and from larger cut-off to smaller, the diminishing effect on the TPE becomes stronger. Figure 5 shows that when using the dipole form factor with $\Lambda = 1.0$ GeV, the effect of the $\hat{V}_{2\pi}^K$ up to about 150 MeV is more or less indistinguishable from the one of the $\hat{V}_{2\pi}^{NPI}$. In other cases, the asymmetry is larger. As illustrated in Fig. 6 the asymmetry using the $\hat{V}_{2\pi}^{DHAL}$ is very sensitive to the used regularizations, because of the form factor modified $\delta$-term, and even exhibits a different sign when the dipole form factor is used. The asymmetry is also chiefly smaller than the reference (NPI $NN$) curve if the regularization is used. In unregularized form, the $\hat{V}_{2\pi}^{DHAL}$ and the $NN$ part of the $\hat{V}_{2\pi}^K$ coincide from which follows that the "w/o" curve in Fig. 6 is identical for each one of these two potentials.
FIG. 6. The same as Fig. 5 but the TPE is by the $\hat{V}^{DHAL}_{2\pi}(r)$ potential and the $NN$ part of the $\hat{V}^{NPI}_{2\pi}(r)$ potential. The "w/o" curve is identical with that of the $NN$ part of the $\hat{V}^{K}_{2\pi}(r)$. This curve is larger than the NPI one, but becomes roughly the same in the case of the $\hat{V}^{K}_{2\pi}$ with dipole form factor and $\Lambda = 1.0$ GeV. Bringing up the difference between the $NN$ and $NN + N\Delta$ for the $\hat{V}^{K}_{2\pi}$ and $\hat{V}^{NPI}_{2\pi}$ compare with Figs. 5 and 1 (or 2,3) respectively. Lastly, minimizing the model dependence of the TPE, one may speculate on the value of the $h^{(1)}_{\pi}$. Some heuristic estimates of it can be "eyeballed" off Figs. 5 and 6 by considering only the "w/o" curves and assuming that the heavy meson effect of Fig. 1 represents realistically the short-range contribution to the analyzing power. The $\hat{V}^{K}_{2\pi}$ suggests that the value of the $h^{(1)}_{\pi}$ should be roughly 50% smaller while $\hat{V}^{DHAL}_{2\pi}$ and $NN$ part of the $\hat{V}^{K}_{2\pi}$ that the "DDH" best value is about correct.

In summary, we have calculated the PNC longitudinal analyzing power $\tilde{A}_L(\vec{p}\bar{p} \rightarrow pp)$ by taking into account the electromagnetic and TPE effects in various models. Coulomb interaction plays virtually no role in the scattering or transmission asymmetries. By using the aforementioned standard set of couplings, we found that the $\hat{V}^{NPI}_{2\pi}$ potential along with the DDH model gives an excellent match with the experimental data. The TPE effect is about two times larger than heavy meson exchange effect throughout the energy scale. Nearly consistent result comes also from the $\hat{V}^{K}_{2\pi}$ potential with the dipole form factor and $\Lambda = 1.0$ GeV cut-off. The above model is also used in the calculation of the cold neutron spin rotation $\frac{d}{dz}\phi$, polarization $\frac{d}{dz}P$, and $\gamma$-asymmetry $A_{\gamma}(\vec{n}p \rightarrow \gamma d)$ in the interaction with parahydrogen $[35]$. Within the said model, the TPE effect reasonably diminishes the OPE
effect by about 10% in the observables. All in all, assuming that the two-pion and heavy meson exchanges are the only major contributions to the analyzing power $\bar{A}_L$, we found that the $\bar{A}_L$ depends mostly on the TPE unless the true value of $h_x^{(1)}$ is significantly smaller than that given by DDH. However, ultimately, the experiments (e.g. the NPDGamma experiment) may decide the reliability of this value. As a conclusion of this work, despite the inescapable model dependence of the observable $\bar{A}_L$, the TPE causes most likely an important effect to it and should not be ignored.

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