Selected topics on Low Energy Antiproton Physics

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Some of the last results on low energy antiproton physics are reviewed. First Faddeev calculations for \(\bar{n}d\) scattering length are presented.

1. Introduction

We review in this contribution some topics on the Low Energy Antiproton Physics that raised our interest since the early shutdown of LEAR, as well as their relation with the theoretical models.

The first section is devoted to \(\bar{N}N\) system. We comment on the \(\bar{n}p\) results obtained by the OBELIX group and discuss with some more details the protonium P-level energy shifts measured by the PS207 experiment.

The second section is devoted to the \(\bar{N}d\) system. We present the first Faddeev calculations for the \(\bar{n}d\) scattering length and review the future perspectives from the theoretical side. Unable to understand the annihilation at a quark level, we have here the possibility to understand it at the level of nucleons.

2. Selected topics on \(\bar{N}N\)

2.1. Low energy \(\bar{n}p\) cross section

The last OBELIX data on \(\bar{n}p\) cross section \[1\], presented in this conference by A. Felicello \[2\], are very astonishing. The structures observed in the total and elastic cross section at antineutron laboratory momenta \(p_{\text{lab}} \approx 70\,\text{MeV}/c\) are not trivial results and can hardly be explained without assuming a nearthreshold state \[3\].

These data have suffered from big fluctuations in the successive steps of their analysis \[4,5\] and the only interesting question – i.e. to what extend they are significant – can not longer be answered at LEAR. The anomalous behaviour indeed concerns only two data points but if they were a consequence of an experimental problem, why should it manifest only in these two intermediate points? As strange as they could appear, these data still remain inside the unitarity constraints. It would be of high interest if the AD could allow low energy scattering experiments. If a direct \(\bar{n}p\) measurement turns out to be not possible, one could see at least whether or not the structure is manifested in the \(\bar{p}p\) cross section.
The problem is modeled by
\[
(E - H_0) \begin{pmatrix} \Psi_{\bar{p}p} \\ \Psi_{\bar{n}p} \\ \Psi_{\bar{n}n} \end{pmatrix} = \hat{V} \Psi \quad \hat{V} = \frac{1}{2} \begin{pmatrix} V_0 + V_1 + 2V_c & 0 & V_0 - V_1 \\ 0 & V_1 & 0 \\ V_0 - V_1 & 0 & V_0 + V_1 \end{pmatrix}
\] (1)
where \( E = T_i + \sum m_i \) denotes the total energy – the same for all channels, \( V_i \) the isospin components of the \( \bar{N}N \) strong potential and \( V_c \) the Coulomb attraction. The \( \bar{n}p \) channel is dynamically decoupled from the two others. However, through the \( T=1 \) component of the \( \bar{N}N \) potential, the \( \bar{n}p \) structure could be visible in the \( \bar{p}p \) elastic cross section near the \( \bar{n}n \) threshold. Indeed, by assuming \( 2\mu_{\bar{p}p} = 2\mu_{\bar{n}p} = 2\mu_{\bar{n}n} = \frac{1}{2}(m_p + m_n) \) the center of mass momenta \( k_i \) of the different channels are given by
\[
k^2_{\bar{p}p} = k^2_{\bar{n}p} + \Delta = k^2_{\bar{n}n} + 2\Delta
\]
where \( \Delta = m_n - m_p = 1.293 \) MeV. The \( \bar{n}n \) threshold \( (k^2_{\bar{n}n}=0) \) correspond to the laboratory moment \( p_{\bar{p}p}=98.6 \) and \( p_{\bar{n}p}=69.7 \) MeV/c, i.e. the region of the observed anomaly. We would like to notice however that this coincidence is only kinematical. Calculations performed in the framework of both optical and unitary coupled-channel models showed that the influence of opening the \( \bar{n}n \) threshold in the \( \bar{p}p \) cross sections indeed exists but it is negligible. If some structure is seen in \( \bar{p}p \) observables near the \( \bar{n}n \) threshold it should have a dynamical origin.

2.2. Protonium P-level shifts

One of the most interesting results from LEAR post-mortem experiments is the \( ^3P_0 \) protonium level shifts obtained in [6]. These authors found the values \( \Delta E_R(3P_0) = -139 \pm 28 \) and \( \Gamma(3P_0) = 120 \pm 25 \) meV. The interest of this result is twofold. First because it is in qualitative agreement with a non trivial prediction of some meson exchange based models. Second for the underlying dynamics that it supports.

Table 1
Protonium P-level shifts \( (\Delta E_R - i\frac{\Gamma}{2} \text{ in meV}) \) and \( \bar{p}p \) scattering volumes \( (a=a_R + i a_I \text{ in fm}^3) \) in some optical models. The \(^3P_0 \) value is compared to experimental result from [3].

| State | \( \Delta E_R \) | \( \frac{\Gamma}{2} \) | \( \Delta E_R \) | \( \frac{\Gamma}{2} \) | \( \Delta E_R \) | \( \frac{\Gamma}{2} \) |
|-------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| \(^1P_1\) | -29. | 13. | -26. | 13. | -24. | 14. |
| \(^3P_0\) | -69. | 48. | -74. | 57. | -62. | 40. |
| \(^3P_1\) | +29. | 11. | +36. | 10. | +36. | 8.8 |
| \(^3P F_2\) | -8.5 | 18. | -4.8 | 15. | -5.9 | 16. |

| & \( a_R \) & \( a_I \) & \( a_R \) & \( a_I \) & \( a_R \) & \( a_I \) |
|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| \(^1P_1\) | -1.19 | -0.53 | -1.07 | -0.52 | -0.99 | -0.58 |
| \(^3P_0\) | -2.81 | -1.99 | -3.01 | -2.31 | -2.53 | -1.62 |
| \(^3P_1\) | +1.22 | -0.47 | +1.46 | -0.42 | +1.48 | -0.36 |
| \(^3P F_2\) | -0.36 | -0.75 | -0.20 | -0.63 | -0.25 | -0.67 |

The predictions from KW [7], DR1 and DR2 [8] optical models for the protonium P-level shifts [9] together with the corresponding \( \bar{p}p \) scattering volumes [10] are given in
Table 1. These models differ from each other both in the annihilation potential and in the meson contents defining its real part. One can see from this table that the $^3P_0$ width is well reproduced and that the energy shift is – in spite of the large predicted values – underestimated by a factor two. We would like to emphasize however that the very prediction of these models is that $^3P_0$ must be a particular case. The reason for $\Delta E_R$ being 2-10 times bigger than other $P$ states is that $V_{^3P_0}$ has an unusually strong and attractive $T=0$ component at distances at which the OBE theory should be reliable (see Figure 1). This large value cannot be reproduced by simple interactions: Bruckner et al. potential, for instance, which proved to be very successful at higher energy [11], gives $a_{sc} = -0.18 - 0.68 i$ fm$^3$. This non trivial result seems be confirmed by experiment and – despite a factor two in $\Delta E_R$ – gives support to meson exchange inspired models.

Figure 1. $V_{NN}(^3P_0)$ in KW model for different values of the cutoff parameter $r_c$. Figure 2. Isospin components for $\bar{N}N$ scattering volumes in KW and DR models.

A little bit more hidden but even more interesting is the way in which the P-wave $\bar{p}p$ scattering volumes are obtained in terms of the $NN$ isospin components. For that purpose we consider the DR1 prediction from Table 1. This value, denoted by $a_{sc} = -3.01 - 2.31 i$, was calculated using $\bar{p}p-\bar{n}n$ coupled equations in which Coulomb ($V_c$) and proton-neutron mass difference ($\Delta$) were included. By removing $\Delta$ and Coulomb corrections, one is left with a purely strong value $a_s$, which in its turn is an average of both isospin components $a_s = \frac{1}{2}(a_0 + a_1)$. The result of this analysis, performed in [10] for all partial amplitudes, is summarized in Table 2.

By applying the same kind of corrections to the experimental scattering volume $a_{pp} = -5.68 - 2.45 i$ fm$^3$ and keeping the same $T=1$ value (this approximation will be justified later) one gets $a_0 \approx -15.7 - 5.77 i$. The large value of its real part is a consequence of the already large measured $\text{Re}(a_{pp}) < 0$ and the destructive interference with the repulsive $T=1$ channel. It cannot be obtained without a nearthreshold $NN$ state which enhances the scattering amplitude. To our opinion, the results of [3] constitutes a direct evidence for its existence. It is interesting to compare this situation with the singlet deuteron channel $a_{np} \approx -20$ fm which inspired the following conclusion: “If deuteron had not been found experimentally one could infer its very existence by the big values of NN scattering
Table 2
Isospin components for the $^3P_0$ scattering volume

| Parameter | DR1               | Exp.                  |
|-----------|-------------------|-----------------------|
| $a_{sc}$  | -3.01 -2.31 i     | (-5.68±1.14) - (2.45±0.53) i |
| $a_s$     | -3.44 -2.66 i     | (-6.5±1.3) - (2.8±0.6) i (*) |
| $a_0$     | -9.58 -5.20 i     | (-15.7±2.6) - (5.8±1.2) i (*) |
| $a_1$     | +2.69 -0.13 i     | +2.69 - 0.13 i (*)     |

(*) obtained by theoretical analysis.

parameters”. Does G-parity manifest itself with such a degree of refinement?

In Figure 2 we have compared the different isospin components of the NN P states as they are given in [10] by the three considered optical models. One can see there, that if theoretical predictions for $^3P_0$ were not ”normal”, the experimental findings at LEAR go well beyond. It is thus interesting to inquire whether or not the existing models are able to reproduce them.

Let us first consider the problem with Im(V)=0 and calculate the $^3P_0$ scattering volume $a$ as a function of the cutoff radius $r_c$. The results for both isospin components are plotted in Figures 3 and 4. The $^3P_0$ state shows a remarkable stability due to the fact that the corresponding potential – displayed in Figure 1 – is repulsive and dominated by the centrifugal barrier. For the $^1P_0$ state, $a$ varies from $-\infty$ to $+\infty$ – what makes easy fitting any experimental result – the vertical asymptotics corresponding to the appearance of new NN bound states. One can also see from these figures that the negative sign of the scattering volume ($a < 0$) implies attractive interaction whereas from $a > 0$ nothing can be inferred.

In presence of annihilation, the results for the $^3P_0$ state are only slightly modified. For a wide range of $r_c$ values, $a$ keeps nearly the same real part and gets only a very small imaginary part, indicating that annihilation cannot come from repulsive channels. On the contrary the behaviour of the $^1P_0$ is much more complex. The pole divergences are reduced to finite oscillations and $a$ can have large imaginary part. In order to agree with
the experimental values, Re(a) and Im(a) should cross the corresponding horizontal line at the same value of $r_c$. If one takes the results from [4] this is not possible even modifying the cutoff parameter in a reasonable range. We have found several solutions to account for that interesting – though disappointing – situation. The first one, see Figure 4, keeps the same annihilation potential but uses a very large value of $r_c$=2 fm. The corresponding potential becomes only of $\sim$100 MeV but we notice that even in this case there is a quasibound state and the experimental data force the model to be in its vicinity. The second solution is obtained by a substantial change in Im(V), thus introducing a quantum number dependence in annihilation potential, as it was the case in Paris potential [12]. Figure 8 shows a solution obtained with a factor 4 reduction in the strength $W_0$ of the imaginary potential ($W_0$=300 MeV instead of 1.2 GeV) and $r_c$ = 0.93 fm.

To conclude this section, we would like to emphasize the importance of low energy experiments as well as the partial wave resolution to get useful information about the interaction. The existence of $^3S_0$ state cannot directly be deduced from the total $pp$
scattering cross section due to its small statistic weight compared to the large number of partial waves involved. In addition, it appears only in the T=0 part of the \( \bar{p}p \) wavefunction and its influence is partially cancelled by the opposite contribution of the T=1 part. It is worth mentioning however that a similar conclusion was also reached in \([3]\) based on an independent analysis.

3. Selected topics on \( \bar{N}d \)

The interest in studying this process – apart from testing \( \bar{N}N \) models in a more complex system – is to have an idea of how \( \bar{N}N \) annihilation takes place in presence of other nucleons. Understanding annihilation at a quark level is made difficult by the fact that we have no appropriate theoretical tools to deal with quarks (q). Calculations are performed on a basis of constituent quarks (Q), mostly in a non relativistic approach. However Q is a very complex object which looks more like \( N \) than q except that it has no asymptotic states.

In \( \bar{N}d \), we know something – at least experimentally – about the elementary amplitudes and it seems in principle easier to inquire about the \( \bar{N}-A \) annihilation mechanism. Does \( \bar{N} \) annihilate on a single \( N \)? If yes, should we expect \( T_{\bar{p}A} \approx T_{\bar{p}N} \)? or rather \( T_{\bar{p}A} \approx T_{\bar{p}N_1} + T_{\bar{p}N_2} \)? and what is the role of the remaining nucleons? Does on the contrary \( \bar{N} \) annihilate on the nuclear bulk? From that purpose deuteron – quite an extended object – is the most interesting nucleus and low energy experiments are the only ones from which one can get useful information. The problem has some analogies with "From where does the e\(^-\) pass through?" in a Young experiment, where we learnt that naive pictures (and questions!) have hard life in Quantum Mechanics.

\( \bar{N}d \) is a genuine problem for the Few-Body Physics community. Faddeev-Yakubovsky (FY) equations in configuration space seem to be the best adapted to deal with non hermitic problems where Coulomb interactions play a vital role. At present, the more complex system for which FY can be exactly solved is \( N=4 \) \([4]\). The only \( \bar{N}A \) problem that can be presently solved is \( \bar{N}-d \). It could be extended in coming years to \( \bar{N}-^3\text{He} \) and – if no bad surprises – to \( \bar{N}-^4\text{He} \). A modest but already fascinating task.

From nuclear physics, we know that the scattering lengths are very dynamical quantities, i.e. very sensible to and determined by the interaction. Any interpretation in terms of geometry is purely lyric. For instance, the singlet n-p value is \( a_{np}^{(1)}=–23 \) fm for a nucleon r.m.s. radius of \( R=0.86 \) fm. On the other hand, the doublet and quartet n-d are respectively \( a_{nd}^{(2)}=0.65 \) fm and \( a_{nd}^{(4)}=6.4 \) fm, i.e. they concern the same geometrical object and differ by one order of magnitude. One should anyway remind that the size of a light nucleus is mainly determined by its binding energy rather than by the number of constituents: deuteron is bigger than \( \alpha \) particle. This sensitivity to details of the interaction made powerful approximate methods fail at very low energies. Only well founded approaches that take into account the full dynamics produce there reliable results.

If the low energy parameters are dynamical, one cannot exclude some \( \bar{N}A \) states having very weak annihilation rates. An example – concerning pd elastic scattering – is given by \( a_{pd}^{(2)} \approx 0 \) fm! In \( \bar{p}p \), one has – according to theoretical predictions \([10]\) – \( \text{Im}(a_{T=0})=0.08 \) fm\(^3\) for \( ^3\text{P}_1 \).

The results that follow have been obtained by solving the Faddeev equations with the
methods developed for \( N=3 \) \cite{15} and \( N=4 \) \cite{14}. In the Faddeev approach, the total \( \bar{n}d \) wave function is obtained as a sum of three amplitudes:

\[
\Psi = \Psi_{pn}(x_1, y_1) + \Psi_{\bar{np}}(x_2, y_2) + \Psi_{\bar{nn}}(x_3, y_3)
\]

(2)
each of them depending on a particular set of Jacobi coordinates:

\[
\begin{align*}
\Psi_p &= \sum_{\alpha_1} \Psi_{i,\alpha_1} \\
\Psi_{\bar{n}} &= \sum_{\alpha_1} \Psi_{i,\alpha_1}
\end{align*}
\]

These amplitudes obey a system of coupled equations

\[
\begin{align*}
(E - H_0 - V_{pn})\Psi_{pn} &= V_{pn}(\Psi_{\bar{np}} + \Psi_{\bar{nn}}) \\
(E - H_0 - V_{\bar{np}})\Psi_{\bar{np}} &= V_{\bar{np}}(\Psi_{\bar{nn}} + \Psi_{pn}) \\
(E - H_0 - V_{\bar{nn}})\Psi_{\bar{nn}} &= V_{\bar{nn}}(\Psi_{pn} + \Psi_{\bar{np}})
\end{align*}
\]

and each of them is in its turn expanded in partial wave components

\[
\Psi_i = \sum_{\alpha_i} \Psi_{i,\alpha_i} \quad \alpha_i = \{l_x, \sigma_x, j_x, l_y, j_y\}
\]

on which the different partial waves of the NN and \( \bar{NN} \) potentials act. Only each separate component has a well defined asymptotic behaviour which, for the scattering problem we are interested in, is:

\[
\begin{align*}
\Psi_{pn}(x_1, y_1) &\approx u_{pn}(x_1) \left[ e^{-iq_1y_1} - S_{11} e^{-iq_1y_1} \right] \\
\Psi_{\bar{np}}(x_2, y_2) &\approx u_{\bar{np}}(x_2) \left[ e^{-iq_2y_2} - S_{12} e^{-iq_2y_2} \right] \\
\Psi_{\bar{nn}}(x_3, y_3) &\approx u_{\bar{nn}}(x_3) \left[ e^{-iq_3y_3} - S_{13} e^{-iq_3y_3} \right]
\end{align*}
\]

where \( u_i(x_i) \) is the \( \bar{NN} \) two-body bound state wavefunction with energy \( E_i = \frac{2a_i^2}{\lambda} \) and \( S_{ij} \) the S-matrix elements. For a given \( \bar{NN} \) model, the values of \( E_i \) depend on the quantum numbers \( \{l_x, \sigma_x, j_x, l_y, j_y\} \) of the \( \bar{NN} \) state.

In practice, all models have one or several \( \bar{NN} \) quasibound states with energies below the deuteron. The spectrum of KW model, for instance, is displayed in Table 3 for both isospin components of the potential as well as for their linear combination \( V_{\bar{nn}} = \frac{1}{2}(V_0 + V_1) \) which governs the \( \bar{n}n \) channel. One can remark that there are no quasibound states for \( \bar{np} \) (\( T=1 \)) but two in \( \bar{nn} \). Even if the initial state has only one asymptotics, e.g. deuteron, the coupled equations will populate all the others and there will be always a non zero probability to leave the initial channel. As a consequence, even in absence of annihilation potential, the \( \bar{n}d \) scattering length will have an imaginary part. We can thus distinguish two kinds of annihilation: the direct one and the one that takes place after the creation of a quasibound \( \bar{NN} \) state. In a practical calculation we cannot avoid taking into account these rearrangement channels.
Table 3
Quasibound states in KW model (MeV)

| State  | $V_0$   | $V_1$  | $V_{\bar{n}n}$ |
|--------|---------|--------|----------------|
| $^3S_0$ | -965-438i | —      | —              |
| $^3SD_1$ | -76-374i  | —      | —              |
| $^3P_0$ | -1155-465i | —      | —              |
| $^3PF_2$ | -619-338i | —      | —              |

The results we have obtained for the two $\bar{n}d$ S-wave states – doublet ($J^\pi = \frac{1}{2}^+$) and quartet ($J^\pi = \frac{3}{2}^+$) – are respectively $a^{(2)}_{\bar{n}d} = 1.59 - 0.88i$ fm and $a^{(4)}_{\bar{n}d} = 1.62 - 0.82i$ fm, what gives an spin-averaged value $\bar{a}_{\bar{n}d} = 1.61 - 0.84i$ fm. Each $\bar{N}N$ Faddeev amplitude involves the $\{l_x, \sigma_x, j_x, l_y, j_y\}$ components listed in Table 4 which include all $V_{\bar{N}N}^{J\leq 1}$ interactions.

It is interesting to compare these results to those given by the $\bar{N}N$ amplitudes which are relevant in $\bar{n}d$. Their spin-averaged values, taken from Table 1 in [10], are $\bar{a}_{\bar{n}n}=0.88-0.84i$ fm and $\bar{a}_{\bar{n}p}=0.85-0.76i$ fm. One then has $\text{Im}(\bar{a}_{\bar{n}d}) \geq \text{Im}(\bar{a}_{\bar{N}N})$, in contrast to the findings of [19] for $\bar{p}d$. In view of this results one is tempted to conclude that $\bar{n}d$ annihilation is dominated by the $\bar{n}n$ channel. However in order to get some light on the annihilation mechanism an analysis of the $\bar{n}d$ wavefunction is required.

Table 4
Faddeev components involved in calculating $\bar{N}d$ scattering lengths

| $J^\pi = 1/2^+$ | $J^\pi = 3/2^+$ |
|-----------------|-----------------|
| $j_x$ | $j_y$ | $\sigma_x$ | $l_x$ | $V$ | $l_y$ | $j_x$ | $j_y$ | $\sigma_x$ | $l_x$ | $V$ | $l_y$ |
| 0    | 1/2  | 0         | 0     | $^1S_0$ | 0     | 0    | 3/2  | 0         | 0     | $^1S_0$ | 2     |
| 1    | 1    | 1         | 1     | $^3P_0$ | 1     | 1    | 1    | 3/2  | 1         | 1     | $^3P_0$ | 1     |
| 1    | 1/2  | 0         | 1     | $^1P_1$ | 1     | 1    | 1/2  | 3/2  | 0         | 1     | $^3S_1$ | 0     |
| 1    | 0    | 1         | 3/2  | $^3P_1$ | 1     | 1    | 1    | 5/2  | 0         | 1     | $^3S_1$ | 2     |
| 1    | 1    | 3/2  | $^3P_1$ | 1     | 1    | 1    | 5/2  | 1     | 1     | $^3P_1$ | 1     |
| 1    | 2    | 3/2  | $^3D_1$ | 1     | 1    | 2    | 5/2  | 2     | 1     | $^3D_1$ | 2     |

We also remark that, contrary to $\bar{n}d$ case, there is very small spin dependence, that is $a^{(2)} \approx a^{(4)}$. This could be expected from the fact that we are dealing with three non identical particles. It is worth noticing however that with $\text{Im}(V)=0$, the values are $a^{(2)} \approx 2.0$ and $a^{(4)} \approx 1.5$. The effect of annihilation potential is to move the two spin
values in opposite directions towards an almost degenerate result. In Figures 9 and 10 is displayed the variation of the real and imaginary parts of $a_{nd}$ as a function of a scaling factor $\lambda$ multiplying $\text{Im}(V)$. They show a smoother behaviour than the one found in the model calculation of [10].

![Figure 9. Evolution of $a_Q$ as a function of the strength of $\text{Im}(V)$.](image9)

![Figure 10. Evolution of $a_D$ as a function of the strength of $\text{Im}(V)$.](image10)

We note also that results can vary up to a factor two depending on the states included in the rearrangement channels. There are however dominant amplitudes which are $^1S_0$ and $^3SD_1$. P-wave seems to account for only $\approx 10\%$ variation.

To our knowledge, there is no any experimental value for $\bar{n}d$ to compare with. Wycech et al. [17] found $a_2 = 1.49 - 0.41i$ and $a_4 = 1.53 - 0.45i$ and even smaller values were found in [18]. In view of comparing these results with those obtained in [19] from $\bar{p}d$ experiment – $\text{Im}(\bar{a})=0.62\pm0.02$ – Coulomb corrections should be included.

These calculations can be improved in several ways. For instance by including D-wave component in deuteron, by coupling the $\bar{nn}$ amplitude to the $\bar{pp}$ one and by disentangling the contribution of the different rearrangement channels. The sensitivity to asymptotic states would also require to use different $\bar{NN}$ models. This work is being pursued by extending to $\bar{Nd}$ P-waves and by including Coulomb interactions in the $\bar{pd}$ system.

4. Conclusions

Some years after its shutdown, LEAR continues to produce interesting results. This shows that its unique characteristics – low energy and high resolution experiments – were very appropriate for understanding the $\bar{NN}$ interaction. If stopping LEAR was a mistake, it would be a fault to stop such an activity, not adapting the AD facility to low energy scattering experiments, against the request of several groups [2,20,21].

We have shown – if needed – the interest in pursuing low energy antiproton physics in three different problems. The intriguing low energy $\bar{np}$ observed structure could be also seen – and thus confirmed – in the elastic $\bar{pp}$ cross section near the charge-exchange threshold. The measured $^3P_0$ protonium level shift constitutes a strong indication of a
NN nearthreshold state in isospin zero channel. This measure should be confirmed and if possible extended to other partial waves. The \( \bar{N}d \) reaction is of interest in understanding the annihilation mechanism. We have presented the first \( \bar{n}d \) results based on exact solutions of Faddeev equations. This theoretical effort should be pursued by extending the calculations to negative parity states – where the \( ^{33}P_0 \) pole could also manifest – and to \( \bar{p}d \) system.

We would like to emphasize that only the knowledge of the isolated partial wave contributions provides effective constraints to models. The \( ^{3}P_0 \) constitutes a nice illustration of how interesting phenomena, not even visible in the cross sections, can manifest after such a separation. For that purpose – and in absence of polarized beams – the atomic experiments seem to be the only ones offering such a possibility.

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