I critically review the status of computations of threshold pion production in proton–proton collisions in the framework of effective field theory approaches or variants thereof. I also present the results of a novel diagrammatic scheme.

1 The problem

Over the last years, very precise data on pion, η and η' meson production in proton-proton collisions in the threshold region have been obtained at IUCF, CELSIUS and COSY (for a review, see the talk by H.O. Meyer1). The basic process is \( pp \rightarrow pNM \), where \( N \) denotes the nucleon and \( M \) a pseudoscalar meson, in our case the \( \pi^0 \), \( \pi^+ \), \( \eta \) or the \( \eta' \). At the respective threshold, the produced meson is soft, i.e. has vanishing three momentum. Consequently, in the case of the pions, which are believed to be the Goldstone bosons related to the spontaneous chiral symmetry breaking QCD is assumed to undergo, this process appears to be a good testing ground for chiral perturbation theory methods. To be specific, at threshold one has only S–waves and thus the pertinent T–matrix for the specific case of \( \pi^0 \) production is parametrized in terms of one single amplitude,

\[
T_{\text{th}}^{\pi^0}(pp \rightarrow pp\pi^0) = \mathcal{A} (i \vec{\sigma}_1 - i \vec{\sigma}_2 + \vec{\sigma}_1 \times \vec{\sigma}_2) \cdot \vec{p}. \tag{1}
\]

The \( \vec{\sigma}_{1,2} \) are the spin–matrices of the two protons. The amplitude is a pseudoscalar quantity and due to the Pauli principle, we are dealing with a \( ^3P_0 \rightarrow ^1S_0 \) s transition (where the 's' refers to the pion angular momentum). The value of the proton cm momentum to produce a neutral pion at rest is given by

\[
|\vec{p}| = \sqrt{M_\pi (m + M_\pi/4)} \simeq \sqrt{M_\pi m} = 362.2 \text{ MeV}, \tag{2}
\]

with \( m = 938.27 \) MeV the proton and \( M_\pi = 134.97 \) MeV the neutral pion mass, respectively. Obviously, \(|\vec{p}|\) vanishes in the chiral limit of zero pion mass. Therefore the soft–pion theorem which requires a vanishing threshold T–matrix in the chiral limit \( M_\pi = 0 \) is trivially fulfilled (as long as \( \mathcal{A} \) does not

---

1 Plenary talk at BARYONS 98, Bonn, Germany, September 1998.
become singular). Stated differently, there is no low–energy theorem for the threshold amplitude $A$. Note also that the value of the threshold momentum is already relatively large. Furthermore, it is known that in the threshold region the strong final–state interactions (FSI) govern the energy dependence. In fact, in the approaches I will discuss in what follows, the complete amplitude is written as

$$T = T^{\text{ISI}} \cdot T^{\text{Prod}} \cdot T^{\text{FSI}},$$

where ISI denotes the initial–state interaction and the microscopic approaches are applied to the production amplitude $T^{\text{Prod}}$. Clearly, such a separation induces a priori some model dependence, as I will discuss later. Let me first describe the existing EFT approaches to calculate/constrain the production amplitude.

## 2 Heavy fermion techniques at work

Most work so far has focused on the production of neutral pions simply because the production amplitude involves, besides many other processes, isoscalar pion–nucleon scattering (see fig.1b), which is known to be suppressed due to chiral symmetry. However, it should be made clear from the start that there must also be important short distance physics, see e.g. and as depicted in fig.1c. The challenge of working out the pion exchange terms in chiral perturbation theory has been been taken up by two groups, one at Seattle and the other at South Carolina. Their calculational schemes are similar. To be definite, I concentrate in what follows on the work of the Seattle group. They use the effective pion–nucleon Lagrangian to second order in small momenta (i.e. tree level graphs), the appearing low–energy constants (LECs) have already been determined from pion–nucleon scattering data. The nucleons are treated as very heavy spin–1/2 fields with four–momentum

$$p_\mu = mv_\mu + k_\mu,$$
with \( v_\mu \) the four–velocity subject to the constraint \( v^2 = 1 \) and \( v \cdot k \ll m \). To leading order, the nucleon propagator is \( S(k) = i/(v \cdot k) \) and one can set up a consistent power counting scheme in terms of small momenta, the pion mass and inverse powers of the nucleon mass. However, their calculations involve a modified counting scheme due to the abovementioned fact that the momentum does not scale like the pion mass, so the heavy fermion propagator for a small residual momentum \( k_\mu \) includes the first kinetic energy correction

\[
\frac{i}{v \cdot k} \rightarrow \frac{i}{k_0 + k^2/2m},
\]

since in the rest frame \( v_\mu = (1, \vec{0}) \). The necessity of such a modification follows from the fact that the heavy baryon formalism can not cope with external momenta as large as \( |\vec{p}| \simeq \sqrt{mM_\pi} \). Let me demonstrate the problems of the heavy baryon approach for a simple (but generic) example. Consider a Feynman diagram which involves a propagating nucleon after emission of the real \( \pi^0 \). I show that this nucleon propagator can not be expanded in powers of \( 1/m \) in the usual way. Let \( v_\mu = (1, \vec{0}) \) be the four–vector which selects the center–of–mass frame. The four–vector of the propagating nucleon is \( mv + k \) with \( k_\mu = (-M_\pi/2, \vec{p}) \) and \( k^2 = -mM_\pi \). Starting on the left hand side with the correct result based on a fully relativistic calculation and then performing the usual \( 1/m \) expansion of the heavy baryon formalism, one has

\[
- \frac{1}{M_\pi} = \frac{1}{v \cdot k + k^2/2m} = \frac{1}{v \cdot k} \sum_{n=0}^{\infty} \left( \frac{-k^2}{2m v \cdot k} \right)^n = - \frac{2}{M_\pi} \sum_{n=0}^{\infty} (-1)^n.
\]

One sees that infinitely many terms of the \( 1/m \)–expansion contribute to the same order. The resulting series does not even converge and oscillates between zero and twice the correct answer. The source of this problem is the extreme kinematics of the reaction \( NN \rightarrow NN\pi \) with \( |\vec{p}| \simeq \sqrt{mM_\pi} \). In that case the leading order operator \( O^{(1)} = iv \cdot \partial \) and the next–to–leading order operator \( O^{(2)} = -\partial \cdot \partial/2m \) lead to the same result, here, \( \pm M_\pi/2 \). This problem is merely related to “trivial” kinematics. Therefore, one either has to modify the propagator as described above or calculate fully relativistically, as discussed in the next section. Space forbids to discuss in big detail the results of these studies, I rather concentrate on the most intriguing ones. First, one finds that the rescattering contribution (fig.1b) comes out with an opposite sign to what one gets in meson–exchange models, i.e. it interferes constructively with the direct production in the Jülich meson–exchange model and destructively in the chiral framework, respectively. Note that while there is still debate about the actual numerical treatment and the ensuing size of the rescattering
contribution in the chiral perturbation theory approaches the sign difference with the meson–exchange model can be considered a genuine feature. This point was addressed in ref. It was argued that the treatment underlying the isoscalar pion–nucleon scattering amplitude and the related transition operator for the process $NN \to NN\pi$ in the chiral framework is not yet sufficiently accurate and thus the resulting rescattering contribution should be considered an artifact of this approximation. Clearly, this does not mean that chiral perturbation theory is invalid but rather that higher order (one loop) effects need to be accounted for. Second, to account for the data, one has to include short distance physics, like e.g. heavy meson exchanges (see fig.1c) and also meson–exchange currents, in particular the anomalous $\pi\rho\omega$ vertex. A typical result is shown in fig.2, which also shows that there is still a sizeable uncertainty related to e.g. the treatment of the isoscalar rescattering.

![Figure 2: Cross-section for $pp \to pp\pi^0$ as function of the pion momentum $\eta$ in units of $M_\pi$ for two NN potentials (Argonne V18 and Reid93) and two parameter sets for the isoscalar $\pi N$ amplitude (ste and cl). Figure courtesy of Bira van Kolck.](image)

The first one–loop calculation (to third order in small momenta) was performed by Gedalin et al. They work in the conventional heavy baryon framework using the lowest order fermion propagator (see the left–hand side of eq.(5)). Some typical one–loop graphs are shown in fig.3. There are, of course, many more diagrams, most of them giving rise to mass and coupling constant renormalization. At this order, there appears also a four–nucleon–pion contact term with a LEC, called $d_1$. This LEC could in principle be fitted from one
total cross section point. The authors of ref.\textsuperscript{10} estimate its values from resonance saturation, specifically by $\rho$ and $\omega$ exchange. This procedure induces, of course, some model–dependence. (The status of resonance saturation in the one–nucleon sector is discussed in ref.\textsuperscript{6}) Therefore, the total cross section can be predicted without any free parameters and it agrees quite nicely with the data in the threshold region. However, there are a couple of loopholes with this result. First, as discussed before, using the leading order nucleon propagator, one can not expect a convergent series. This is reflected in the results of ref.\textsuperscript{10}:

$$T_{\text{rescattering}} : T^{\text{loop}} : T^{\text{direct}} \simeq 5 : 2 : 1,$$

(7)

according to figs.1b,3,1a, respectively. Without the loop contribution, the total cross section is underestimated by a factor of about 2.5. One can pin down this problem of the bad convergence more specifically. The class of loop graphs shown in fig.3 is directly proportional to the scalar form factor of the nucleon at $t = M_{\pi}^2 \simeq 0.1 \text{GeV}^2$. In the heavy fermion approach, it leads to a very large enhancement of the production amplitude. If one, however, calculates these loop diagrams fully relativistically and evaluates the scalar form factor, its contribution is very small.\textsuperscript{12} This simply reflects the fact that a one–loop calculation at third order is not sufficiently accurate, even the shift of the scalar form factor to the Cheng–Dashen point at $t = 4M_{\pi}^2 \simeq 0.02 \text{GeV}^2$ is off by a factor of two. Stated differently, the relativistic calculation shows that higher order $1/m$ terms are large. Similar remarks apply to the full isoscalar pion–nucleon scattering amplitude. From these results I conclude that heavy baryon chiral perturbation theory is presumably not the appropriate framework to gain a deeper understanding of the role of chiral dynamics in the process $pp \rightarrow pp\pi^0$. As already pointed out in ref.\textsuperscript{10}, a more detailed study of the reaction $pp \rightarrow d\pi^+$ could, however, pave the way to a deeper understanding of the pion dominated part of the transition operator (simply because short–distance physics plays a much smaller role in this process). Finally, I remark that the treatment discussed so far involves an unavoidable ambiguity due to the exchanged pion (say in the rescattering graph, fig.1b) being off–shell. For a given prescription of defining the pion field, one can perform an off–shell extrapolation but this
will always be model–dependent – in quantum field theory only on–shell matrix elements and transition currents can be calculated.

3 Diagrammatic approach

Based on the observation that the heavy baryon framework is not expected to converge, in ref. a different approach was pursued. Consider first the process \( pp \rightarrow pp\pi^0 \). Approximating the near threshold T–matrix by the T–matrix exactly at threshold one gets for the unpolarized total cross section

\[
\sigma_{\text{tot}}(T_{\text{lab}}) = |A|^2 \int dW \, KF(W) \, F_p(W),
\]

where the flux and three-body phase space factors, denoted as \( KF(W) \), can be approximated by an analytical expression which is accurate within a few percent in the threshold region. \( F_p(W) \) is the correction factor due the final–state interaction, with \( W \) the final–state invariant di–proton mass. It is evaluated in the effective range approximation,

\[
F_p(W) = \left\{ 1 + \frac{a_p}{4}(a_p + r_p)P^2_s + \frac{a_p^2 r_p^2}{64} P^4_s \right\}^{-1},
\]

with \( P^2_s = W^2 - 4m^2 \) and \( a_p \) (\( r_p \)) the \( pp \) scattering length (effective range) including electromagnetic corrections. This is of course a very strong assumption but it allows one to explain the energy dependence of the experimental total cross sections very accurately in terms of a single constant amplitude \( A \). At the end of this section, I derive this particular treatment of the FSI from an effective field theory (EFT) approach. Separating off the final–state interaction in that way, one can then pursue a diagrammatic approach to the (on-shell) production amplitude \( A \). This allows one to investigate in a simple fashion the role of one-pion exchange and chiral loop effects together with shorter range exchanges due to heavier mesons. In a similar fashion, one can investigate the other processes \( pp \rightarrow pn\pi^+ \), \( pp \rightarrow ppp \) and \( pp \rightarrow pp\eta' \) (see below). To be specific, consider the \( pp\pi^0 \) reaction. Assuming the form eq. (8), the S–wave threshold amplitude can be extracted from the data,

\[
A^{(\text{emp})} = (2.7 - i 0.3) \text{ fm}^4,
\]

as shown by the solid line in fig.4. The imaginary part is due to the \( ^3P_0 \) \( pp \) phase shift taken at the threshold energy in the lab frame, \( T_{\text{th}}^{\text{lab}} = 279.65 \text{ MeV} \), where \( \delta^{(3P_0)} = -6.3^\circ \). Thus the imaginary part \( \text{Im} \, A \) is about \(-1/9\) of the real part \( \text{Re} \, A \) and contributes negligibly to the total cross section near threshold. This
number can be well understood in terms of chiral $\pi^0$ exchange (including chiral $\pi^0$ rescattering) and heavy meson ($\omega$, $\rho$, $\eta$) exchanges based on a relativistic Feynman diagram calculation with all parameters being fixed from reliable methods, like forward NN dispersion relations for the $\omega$ and a dispersion–theoretical analysis of $\bar{N}N \rightarrow \pi\pi$ for the $\rho$. Interestingly, in the relativistic approach the rescattering contribution interferes constructively with the direct production term (cf. fig.1a), in contrast with the heavy baryon approach and in agreement with the meson–exchange models. One can also evaluate some classes of one-loop graphs and finds that they lead to small corrections of the order of a few percent. Therefore chiral loops do not seem to play any significant role in the processes $NN \rightarrow NN\pi$, which are dominated by one–pion exchange and short–range physics. I remark also that both the long range $\pi^0$ exchange and the short range vector meson exchange lead to contributions to the threshold amplitude $A$ which do not vanish in the chiral limit $M_\pi \rightarrow 0$. There is no chiral suppression of the reaction $pp \rightarrow pp\pi^0$ compared to other $NN\pi$ channels. In all cases the respective threshold amplitudes are non-zero (and finite) in the chiral limit. This is in contrast to the widespread believe that $pp \rightarrow pp\pi^0$ is suppressed for reasons of chiral symmetry.

Within the same approach, one can investigate the threshold behavior of the process $pp \rightarrow pn\pi^+$. It is given in terms of $A$ and the triplet threshold...
amplitude $\mathcal{B}$,

$$T_{\text{th}}^{\eta}(pp \rightarrow pn\pi^+) = \frac{A}{\sqrt{2}} (i \vec{\sigma}_1 - i \vec{\sigma}_2 + \vec{\sigma}_1 \times \vec{\sigma}_2) \cdot \vec{p} - \sqrt{2} B i(\vec{\sigma}_1 + \vec{\sigma}_2) \cdot \vec{p}. \quad (11)$$

Here, the second term refers to the $^3P_1 \rightarrow ^3S_1$ transition. Note also that the singlet transition is suppressed because of the large singlet scattering length. From the IUCF data, one can determine the empirical value of the triplet amplitude, $\mathcal{B}^{\text{emp}} = (2.8 - i 1.5) \text{ fm}^4$. The large imaginary part is related to the strong ISI in the $^3P_1$ entrance channel, which naturally cannot be explained by tree graphs only. The value of the $^3P_1$ phase at the threshold energy for $pp \rightarrow pn\pi^+$ is $-28.1^\circ$. The corresponding real part $\text{Re } \mathcal{B}$ is well reproduced by chiral one–pion exchange and short–range vector meson physics. For the same parameters as used in the study of $pp\pi^0$, one gets $\text{Re } \mathcal{B} = 2.74 \text{ fm}^4$. A more detailed discussion of this channel and the problems related to the large imaginary part can be found in ref. 12.

Let me now turn to $\eta$ production. The threshold matrix element takes the form

$$T_{\text{th}}^{\eta}(pp \rightarrow pp\eta) = C (i \vec{\sigma}_1 - i \vec{\sigma}_2 + \vec{\sigma}_1 \times \vec{\sigma}_2) \cdot \vec{p}. \quad (12)$$

with $C$ the (complex) threshold amplitude for $\eta$–production. The $\eta$–production threshold is reached at a proton laboratory kinetic energy $T_{\text{lab}}^{\text{th}} = M_\eta(2 + M_\eta/2m) = 1254.6 \text{ MeV}$, where $M_\eta = 547.45 \text{ MeV}$ denotes the eta–meson mass. In the case of $\eta$–production near threshold it is also important to take into account the $\eta p$ final–state interaction, since the $\eta N$–system interacts rather strongly near threshold. In fact a recent coupled–channel analysis of the $(\pi N, \eta N)$–system finds for the real part of the $\eta N$ scattering length $\text{Re } a_{\eta N} = (0.717 \pm 0.030) \text{ fm}$. For comparison, this value is about six times larger than the $\pi^–p$ scattering length, $a_{\pi^–p} = 0.125 \text{ fm}$, as measured e.g. in pionic hydrogen. In ref. 13, it is assumed that the correction due to the $S$–wave $\eta p$ FSI near threshold can be treated in effective range approximation analogous to the $S$–wave $pp$ FSI. The further assumption is made that the FSI in the $pp$ subsystem and in the two $\eta p$ subsystems do not influence each other and that they factorize. The corresponding form of the unpolarized total cross section in terms of the $S$–wave amplitude $C$ the various FSI functions $F_p(W), F_\eta(s_\eta)$ can be found in ref. 14. Note that the $\eta p$ FSI function is complex–valued (because the $\eta N$ scattering length has a sizeable imaginary part). From the CELSIUS data, one finds for the modulus of the threshold amplitude $|C| = 1.32 \text{ fm}^4$. The resulting energy dependent cross section from threshold up to $T_{\text{lab}} = 1375 \text{ MeV}$ is shown in fig. 6 together with the data from CELSIUS. It is rather astonishing that one can describe the total cross section
Figure 5: The eta–production cross section $\sigma_{\text{tot}}(pp \rightarrow pp\eta)$ as a function of $T_{\text{lab}}$. The data are taken from CELSIUS.

data up to 100 MeV above threshold with a constant threshold amplitude $C$ and a simple factorization ansatz for the three–body FSI. The relativistic Feynman graphs contributing to the threshold amplitude $C$ can be easily evaluated. To account for the strong $\eta N$ S-wave rescattering, which is often attributed to the nucleon resonance $S_{11}(1535)$, one has to introduce a local $NN\eta\eta$ contact term, $L_{\eta N} = K\bar{N}(x)N(x)\eta^2(x)$. Its strength $K$ can be obtained from fitting the scattering length $a_{\eta N}$. With $g_{\eta N} = 5.3$, which is close to flavor SU(3) estimates and determinations from boson–exchange models, one can exactly reproduce the value of $C$. It is worth pointing out that $\rho^0$ exchange is the dominant contribution, because it is enhanced by factors of $M_\eta/m \simeq 0.6$ quite in contrast to the neutral pion case, where $\omega$ and $\pi$ exchange are the dominant mechanisms. One can also predict the threshold S–wave amplitude $D$ for the process $pn \rightarrow pn\eta$. One finds $|D| = 1.8 \text{fm}^4$, somewhat below the experimental value, $|D|^{\text{emp}} = 2.3 \text{fm}^4$, as determined from the recent data on quasifree production off the deuteron measured at CELSIUS.

At COSY, $\eta'$ production very close to threshold has been measured for the first time, at much smaller excess energies than the few data points from the now defunct SATURNE facility. These data seem to follow three–body phase space with little indication of FSI. Note, however, that so close to threshold the energy determination is difficult and thus there is a sizeable uncertainty in the value of the excess energy (for more details, see ref.14). Within the diagrammatic approach, the empirical S–wave amplitude can be obtained with

$$g_{\eta' NN} (1 - 1.28\varepsilon) = 1.12,$$  \hspace{1cm} (13)
with \( g_{\eta'NN} \) the so far undetermined \( \eta' - N \) coupling constant and \( \varepsilon \) the pseudoscalar to pseudovector mixing parameter. Since the \( \eta' \) is not a Goldstone boson, one is not forced to use only derivative couplings. Interestingly, only the tensor interaction of the \( \rho \)-exchange (\( \sim \kappa_{\rho} \)) is sensitive to the parameter \( \varepsilon \). Combining this result with constraints from deep inelastic lepton–nucleon scattering (the spin content of the nucleon), one infers

\[
g_{\eta'NN} = 2.5 \pm 0.7 \, , \quad \varepsilon = 0.4 \pm 0.1 \, .
\]

(14)

It remains to be seen whether other \( \eta' \)-production processes (e.g. photoproduction \( \gamma p \to \eta' p \)) are consistent with these values.

Finally, I want to give an elementary derivation of the final-state interaction correction factor \( F_p(W) = \left[ 1 + \frac{a_p^2 P^2_*}{1 + i a_p P_*} \right]^{-1} \), \( P^2_* = W^2/4 - m^2 \), in the scattering length approximation (i.e. for the effective range \( r_p = 0 \)). Close to threshold all final state three–momenta are small and therefore one can approximate both the meson production process \( NN \to NN\pi^0 \) and elastic \( NN \to NN \) scattering by momentum independent contact vertices proportional to \( A \) and the scattering length \( a_p \), respectively. Consider first low energy \( NN \)-scattering in this approximation. The bubble diagrams with 0,1,2,\ldots rescatterings can be easily summed up in the form of a geometric series,

\[
a_p - 4\pi a_p^2 \int \frac{d^3 l}{(2\pi)^3} \frac{1}{P^2_* + i0 + l^2 + \cdots} = a_p + i a_p^2 P_* + \cdots = \frac{a_p}{1 - i a_p P_*} ,
\]

(15)

using dimensional regularization (in the \( \overline{MS} \)-scheme) to evaluate the (vanishing) real part of the loop integral. Obviously, the sum of these infinitely many loop diagrams is just the unitarized scattering length approximation which leads to \( \tan \delta_0(W) = a_p P_* \). Next, consider in the same approximation meson production followed by an arbitrary number of \( NN \)-rescatterings in the final–state. Again, these loop diagrams can be summed up to,

\[
\frac{A}{1 - i a_p P_*} ,
\]

(16)

and taking the absolute square,

\[
\left| \frac{A}{1 - i a_p P_*} \right|^2 = \frac{|A|^2}{1 + a_p^2 P^2_*} ,
\]

(17)

one encounters the final–state interaction correction factor \( F_p(W) = \left[ 1 + a_p^2(W^2/4 - m^2) \right]^{-1} \) in scattering length approximation. Since the scattering length is much bigger than the effective range parameter for \( NN \)-scattering.
(a_p >> r_p) one has already derived the dominant effect due to the FSI. Of course, in order to be more accurate one should eventually go beyond momentum independent contact vertices. The main point, however, I want to make here is that the FSI correction factor $F_p(W)$ (for $r_p = 0$) has a sound foundation in effective field theory.

Of course, this approach also has some drawbacks. First, in view of the recent polarization measurements at IUCF and RCNP, one should go beyond the S–wave approximation. Second, it is not obvious how to systematically improve the calculations and third, the treatment of the FSI might be too simplistic (for a recent discussion, see e.g. ref. 13).

4 Outlook

I given a short review of EFT or EFT–inspired approaches to pion (as well as $\eta$ and $\eta'$) production in proton–proton collisions. The much discussed reaction $pp \to ppp^0$ does not seem to offer a testing ground for chiral dynamics due to large uncertainties related to chiral pion–exchange, the short distance physics and the strong ISI/FSI. Even the novel EFT scheme due to Kaplan, Savage and Wise 17 is (in its present formulation) not able to cope with pion momenta as large as in this reaction. On the other hand, there exist also the rather successful meson–exchange models based on effective Lagrangians from the Jülich (and RCNP) group(s). While these work fairly well, there is no systematics means of improving them and also, they employ highly model–dependent meson–nucleon form factors, which have no sound physical basis. Since a consistent EFT approach to deal with these processes does not seem to be on the horizon, it might be worthwhile to try to constrain the meson–exchange models with EFT ideas and eventually gain a better understanding of the meson–nucleon interaction regions. Also, more effort should be invested in studying the process $pp \to d\pi^+$ since it is much less contaminated by short–distance physics. Theory still has a long way to go to catch up with the tremendous precision obtained in the recent experiments.

Acknowledgments

I thank Dan-Olof Riska for giving me the opportunity to present these thoughts and Silas Beane for a critical reading of the manuscript. I also thank the Institute for Nuclear Theory at the University of Washington for its hospitality and the Department of Energy for partial support during the completion of this work.
References

1. H.O. Meyer, these proceedings.
2. G. Miller and P. Sauer, Phys. Rev. C 44, 1725 (1991).
3. T.-S.H. Lee and D.O. Riska, Phys. Rev. Lett. 70, 2237 (1993).
4. T.D. Cohen, J.L. Friar, G.A. Miller and U. van Kolck, Phys. Rev. C 53, 2661 (1996); U. van Kolck, G.A. Miller and D.O. Riska, Phys. Lett. B 388, 679 (1996).
5. B.Y. Park, F. Myhrer, J.R. Morones, T. Meissner and K. Kubodera, Phys. Rev. C 53, 1519 (1996).
6. V. Bernard, N. Kaiser and Ulf-G. Meißner, Nucl. Phys. A 615, 483 (1997); Nucl. Phys. B 475, 147 (1995).
7. C. Hanhart, J. Haidenbauer, A. Reuber, C. Schütz and J. Speth, Phys. Lett. B 358, 21 (1995); C. Hanhart et al., nucl-th/9808020.
8. T. Sato, T.-S.H. Lee, F. Myhrer and K. Kubodera, Phys. Rev. C 56, 1246 (1997).
9. C. Hanhart, J. Haidenbauer, M. Hoffmann, Ulf-G. Meißner and J. Speth, Phys. Lett. B 424, 8 (1998).
10. E. Gedalin, A. Moulem and L. Razdolskaya, nucl-th/9803023.
11. V. Bernard, N. Kaiser and Ulf-G. Meißner, Int. J. Mod. Phys. E 4, 193 (1995).
12. V. Bernard, N. Kaiser and Ulf-G. Meißner, nucl-th/9806013.
13. M. Batinic, I. Dadic, I. Slaus, A. Svarc, B.M.K. Nefkens and T.-S.H. Lee, Physica Scripta (1998) in print.
14. P. Moskal et al., Phys. Rev. Lett. 80, 3202 (1998).
15. A.V. Efremov, J. Soffer and N.A. Törnqvist, Phys. Rev. D 44, 1369 (1991); G. Altarelli, R.D. Ball, S. Forte and G. Ridolfi, Nucl. Phys. B 496, 337 (1997).
16. C. Hanhart and K. Nakayama, nucl-th/9809059.
17. D.B. Kaplan, these proceedings.