Kaonic atoms and in-medium $K^-N$ amplitudes

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Abstract

Recent work on the connection between in-medium subthreshold $K^-N$ amplitudes and kaonic atom potentials is updated by using a next to leading order chirally motivated coupled channel separable interaction model that reproduces $\bar{K}N$ observables at low energies, including the very recent SID-DHARTA results for the atomic $K^-\text{-hydrogen}$ 1s level shift and width. The corresponding $K^-$-nucleus potential is evaluated self-consistently within a single-nucleon approach and is critically reviewed with respect to empirical features of phenomenological optical potentials. The need to supplement the single-nucleon based approach with multi-nucleon interactions is demonstrated by showing that additional empirical absorptive and dispersive terms, beyond the reach of chirally motivated $K^-$-nucleus potentials, are required in order to achieve good agreement with the bulk of the data on kaonic atoms.

Keywords: in-medium subthreshold scattering amplitudes, coupled channel chiral models, kaonic atoms

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1. Introduction

The bulk of the data on strong interaction effects in kaonic atoms is due to experiments of over three decades ago. With the exception of the very light atoms of $K^-\text{H}$ and $K^-\text{He}$ the rest of the data could be described rather well with the help of $K^-$-nucleus optical potentials $[1,2]$. Recent experiments on $K^-\text{H}$ and $K^-\text{He}$ with much reduced background removed the ‘puzzles’ with these two atoms $[3,4,5,6]$. However, the depth of the attractive $K^-$-nucleus real potential, which in phenomenological analyses came out in the range of 150-200 MeV $[7]$, presented a theoretical challenge in as much as in-medium chiral threshold $K^-N$ scattering amplitude input led to a lower value of order 120 MeV $[8]$, or even to a considerably lower value of 40-50 MeV $[9]$. 

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This outstanding discrepancy is of current interest, since it is relevant to the role of $K^-$ mesons in multistrange self-bound matter and in compact stars \cite{10, 11}. The problem has been largely resolved very recently \cite{12, 13} noting that in-medium chiral \textit{subthreshold} $K^-N$ scattering amplitudes provide the relevant input and thereby demonstrating the need to supplement the model by multi-nucleon terms, as discussed in the present work. This leads to deep real potentials in agreement with the purely phenomenological analyses.

The present paper is an update of Refs. \cite{12, 13}, based on a recent in-medium coupled channel chirally motivated separable interaction model which produces good fits to all the low energy antikaon-nucleon data, including the latest $K^-H$ atom results from the SIDDHARTA experiment \cite{5}. Section 2 outlines the self-consistent handling of \textit{subthreshold} $K^-N$ amplitudes while section 3 deals with the resulting $K^-$-nucleus amplitudes. Section 4 reports on global optical model fits to kaonic atom data and the last section provides summary and conclusions.

\section{2. $K^-N$ scattering amplitudes}

The potential experienced by a $K^-$ meson of energy $E_{K^\text{lab}} = \omega_K$ interacting with a nucleus of density $\rho$ is given in the single-nucleon approximation by

$$V_{K^-}(\omega_K; \rho) = -\frac{2\pi}{\omega_K} (1 + \frac{\omega_K}{m_N}) F_{K^-N}(\vec{p}, \sqrt{s}; \rho) \rho,$$

where $F_{K^-N}(\vec{p}, \sqrt{s}; \rho)$ is the in-medium $K^-N$ scattering amplitude, reducing in the low-density limit $\rho \to 0$ to the free-space $K^-N$ c.m. forward scattering amplitude $F_{K^-N}(\vec{p}, \sqrt{s})$, $\vec{p}$ is the relative $K^-N$ momentum, $s = (E_K+E_N)^2 - (\vec{p}_K + \vec{p}_N)^2$ is the Lorentz invariant Mandelstam variable equal to the square of the total $K^-N$ energy in the two-body c.m. frame, and the nucleon energy $E_N$ is approximated by its mass $m_N$ in the kinematical factor in front of $F_{K^-N}$. The in-medium amplitude $F_{K^-N}(\vec{p}, \sqrt{s}; \rho)$ in this work is a chirally motivated amplitude constructed within a full octet $0^-$ meson--octet $1/2^+$ baryon coupled channel separable interaction model \cite{14, 15} which in its latest next to leading order (NLO) version NLO30 \cite{15} incorporates the recent SIDDHARTA data for the atomic $K^-H 1s$ level shift and width \cite{5}. In this separable interaction model, the in-medium coupled channel scattering amplitudes assume the form

$$F_{ij}(p, p'; \sqrt{s}; \rho) = g_i(p) f_{ij}(\sqrt{s}; \rho) g_j(p'),$$
with form factors $g_j(p) = \alpha_j^2/(p^2 + \alpha_j^2)$. The momentum dependence intro-
duced by the form factor $g_{K^-N}$ in the separable interaction model of
Refs. [14, 15] is relatively weak for the applications discussed in the present
work and is secondary to the strong energy dependence of the reduced am-
plitude $f_{K^-N}$ generated by the $\Lambda(1405)$ subthreshold resonance.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure1.pdf}
\caption{Energy dependence of the c.m. $K^-N$ reduced amplitude $(3)$ in version NLO30
of the chiral model [15] below and above $E_{th} = m_K + m_N = 1432$ MeV. Dashed curves:
free-space amplitude; dot-dashed curves: Pauli blocked amplitude at $0.5 \rho_0$; solid curves:
including meson and baryon self energies (SE), also at $0.5 \rho_0$.}
\end{figure}

Free-space and in-medium reduced amplitudes for half nuclear matter
density $\rho = 0.5 \rho_0$ are shown in Fig. 1 for the isospin-averaged combination

$$f_{K^-N} = \frac{1}{2}(f_{K^-p} + f_{K^-n}) = \frac{3}{4}f_{I=1} + \frac{1}{4}f_{I=0},$$

(3)
corresponding to symmetric nuclear matter. Similar results are obtained at
full nuclear matter density $\rho_0 = 0.17$ fm$^{-3}$. Of the two in-medium amplitudes
shown in the figure, the one marked “without SE” imposes Pauli blocking on
intermediate $\bar{K}N$ states for $\rho \neq 0$ [8], whereas the one marked “with SE” adds
self consistently hadron self energies in intermediate states [16], following a
procedure suggested in Ref. [17]. The real part of all three amplitudes exhibits strong energy dependence, switching from weak attraction above \( K^-N \) threshold to strong attraction below threshold. As a rule of thumb, \( \text{Re } f = 1 \) fm translates into a sizable attraction \( \text{Re } V_{K^-} \approx -100 \) MeV. The imaginary part of these amplitudes exhibits a peak, related to the subthreshold \( \Lambda(1405) \) resonance, with a steep decrease at lower energies, becoming vanishingly small near the \( \pi\Sigma \) threshold about 100 MeV below the \( K^-N \) threshold. In between the two limits of the energy scale in Fig. 1, \( E - E_{\text{th}} = \pm 100 \) MeV, the three amplitudes differ appreciably from each other. At threshold, in particular, the real part of the “with SE” amplitude is about half of that “without SE”, corresponding to a depth \(-\text{Re } V_{K^-}(\rho_0) \approx -40 - 50 \) MeV, in agreement with Ramos and Oset [9].

It was recognized in the early 1970s that the strong energy dependence of the two-body amplitude, particularly in the subthreshold region where the \( \bar{K}N \) quasibound state \( \Lambda(1405) \) dominates, provides the underlying structure for \( K^- \) nuclear interactions at and near the \( K^- \) nucleus threshold \([18, 19, 20]\). This idea has been reformulated and applied recently in Refs. [12, 13] to a comprehensive study of kaonic atoms. The essential idea is to replace the two-body variables \( \vec{p} \) and \( \sqrt{s} \) of the in-medium scattering amplitude by appropriate density dependent averages in the nuclear medium. This may be summarized by the following relationships:

\[
\sqrt{s} \rightarrow E_{\text{th}} - B_N - B_K - \xi_N \frac{p_N^2}{2m_N} - \xi_K \frac{p_K^2}{2m_K}, \tag{4}
\]

upon neglecting quadratic terms in the binding energies \( B_K = m_K - E_K, B_N = m_N - E_N \) near threshold \( (E_{\text{th}} = m_N + m_K) \), and

\[
p^2, \ p'^2 \rightarrow \xi_N \xi_K (2m_K \frac{p_N^2}{2m_N} + 2m_N \frac{p_K^2}{2m_K}), \tag{5}
\]

where \( \xi_{N(K)} = m_{N(K)}/(m_N + m_K) \) in both of these substitutions. Replacing in Eqs. (4) and (5) the kinetic energy \( p_K^2/(2m_K) \) in the local density approximation by \(-B_K - \text{Re } V_{K^-}(\rho)\) where \( V_{K^-} = V_{K^-} + V_c \), with \( V_c \) the \( K^- \) Coulomb potential generated by the finite-size nuclear charge distribution, and approximating the nucleon kinetic energy \( p_N^2/(2m_N) \) in the Fermi gas model by \( 23 (\rho/\rho_0)^{2/3} \) MeV, Eqs. (4) and (5) become

\[
\sqrt{s} \approx E_{\text{th}} - B_N - \xi_N B_K - 15.1(\frac{\rho}{\rho_0})^{2/3} + \xi_K \text{Re } V_{K^-}(\rho), \quad \text{(in MeV)} \tag{6}
\]
where all the terms following $E_{\text{th}}$ on the r.h.s. are negative, thus implementing the anticipated downward energy shift into the $K^-N$ subthreshold energy region, and

$$p^2 \approx \xi_N \xi_K [2m_K 23 (\rho/\rho_0)^{2/3} - 2m_N (B_K + \text{Re} V_K^- (\rho))], \quad \text{(in MeV)}$$

(7)

where both terms on the r.h.s. are positive for attractive potentials $V_K^-$. The dominant contribution in Eqs. (6) and (7) arises from $\text{Re} V_K^- (\rho)$, resulting in a downward energy shift of up to 60 MeV as demonstrated in Fig. 2 and in values of $p(\rho_0), p'(\rho_0)$ as high as 275 MeV/c for $K^-$ nuclear potential depths reaching 180 MeV in phenomenological studies [1]. These momenta are well within the NLO30 momentum dependence scale $\alpha_{KN} = 700$ MeV/c.

![Graph showing subthreshold energies as function of nuclear density.](image)

Figure 2: Subthreshold energies as function of nuclear density, see text.

Having transformed the dependence of the in-medium scattering amplitudes $F_{K^-j}(\vec{p}, \sqrt{s}, \rho)$ ($j = p, n, N$) on $\vec{p}$ and $\sqrt{s}$ into a density dependence, we denote the resultant in-medium scattering amplitudes by $F_{K^-j}(\rho)$. In order to allow for different proton and neutron distributions in the actual calculations detailed below, the in-medium amplitude $F_{K^-N}(\rho)$ which should substitute for $F_{K^-N}(\vec{p}, \sqrt{s}, \rho)$ in the construction of the $K^-$-nucleus potential Eq. (1) is further replaced by an effective in-medium amplitude $F_{K^-N}^{\text{eff}}(\rho)$:

$$F_{K^-N}^{\text{eff}}(\rho)(r) = F_{K^-p}(\rho)p_p(r) + F_{K^-n}(\rho)p_n(r),$$

(8)
with $\rho_p$ and $\rho_n$ normalized to $Z$ and $N$, respectively, and $Z + N = A$. The reduced amplitudes $f_{K^-p}$ and $f_{K^-n}$ are evaluated at $\sqrt{s}$ given by Eq. (6), where the $K^-$ atomic binding energy $B_K$ is neglected with respect to the average nucleon binding energy $B_N \approx 8.5$ MeV. A similar approximation is made in Eq. (7) for $p^2$ when using the form factors $g_{K^-N}(p)$ of Eq. (2).

The $K^-$-nucleus potential $V_{K^-} (\rho)$ is calculated by requiring self consistency in solving Eq. (6) with respect to Re $V_{K^-}$, i.e., the value of Re $V_{K^-}(\rho)$ in the expression for $\sqrt{s}$ and in the form factors $g_{K^-N}$ has to agree with the resulting Re $V_{K^-}(\rho)$. This is done at each radial point and for every target nucleus in the data base.

3. $K^-$-nucleus scattering amplitudes

The present model transforms the energy dependence of subthreshold effective amplitudes, Eq. (8), into density dependence. This transformation is hardly sensitive to the nucleus involved, as is seen in Fig. 2 calculated in the “without SE” version of the NLO30 model. In the “with SE” version the energies for a given density differ from the plotted values by 2-3 MeV. The energy shifts do not vanish for zero density because we have used a fixed average nucleon binding energy of 8.5 MeV. Replacing it by a position-dependent $B_N \rightarrow B_N \rho(r)/\bar{\rho}$, in order to satisfy the low-density limit, causes the energy shift to vanish far outside the nucleus, with minor overall effects on the present results. It is seen from the figure that, e.g., for a density of 50% of nuclear matter density, the downward energy shift is $\approx 40$ MeV which, from Fig. 1, implies a real amplitude at least twice larger than the threshold value in the “with SE” version.

The density dependencies of the real and of the imaginary part of the effective amplitude $F_{K^-N}^{\text{eff}}$ are of particular interest because they are related to characteristic features of phenomenological optical potentials. It was shown already in 1993 \cite{7} that with empirical density-dependent potentials, where the effective $K^-N$ scattering amplitude within a $t\rho$ model depends on the density, improved fits to the data were obtained compared to fits using fixed amplitudes. It was observed that in addition to the increased depth of the

\footnote{The precise value used for $B_N$ in our kaonic atom global fit hardly matters within reasonable limits. Sensitivity to $B_N$ is expected in studies limited to light nuclei, as exhibited recently by analyzing FINUDA data \cite{21} of $\Lambda$ hypernuclear formation with stopped $K^-$ mesons on targets from lithium to oxygen \cite{22}.}
best-fit real potentials, these were characterised by *compression* relative to the corresponding nuclear densities, with r.m.s. radii of the real potential smaller than the corresponding nuclear radii. Reduced r.m.s. radii of optical potentials mean that the underlying in-medium $K^-N$ interaction increases with density, a robust feature that is insensitive to details.

![Graph showing effective amplitudes as function of nuclear density in model NLO30 without SE.](image)

**Figure 3:** Effective amplitudes as function of nuclear density in model NLO30 without SE.

Figure 3 shows the NLO30 effective amplitudes for Ni and Pb as function of nuclear density, calculated in the “without SE” version. Qualitatively similar results are obtained also in the “with SE” version, see Ref. [12]. Regions of low density, i.e. large radii, are the most effective in determining the r.m.s. radius of a distribution and the sharp rise of the real part in the extreme surface region can lead to compression of the real potential. The opposite dependence is observed for the imaginary part, thus implying *inflation* of the imaginary potential relative to the nuclear density. Quantitatively, however, the change of r.m.s. radii relative to nuclear densities is found to disagree with the empirical trends of Ref. [7], with too little compression for the real part and far too strong inflation for the imaginary part. Inevitably this is reflected in the quality of agreement with the data, as demonstrated in the next section.

Strong-interaction effects in kaonic atoms are dominated by absorption, as is evident from level widths being significantly larger than the correspond-
ing level shifts. Moreover, the shifts are always repulsive although the real potential is attractive, again pointing to the dominance of absorption. Therefore it is argued that the above marked decrease of the imaginary part of the effective scattering amplitude is the main deficiency of the present model, a decrease originating in the sharp decrease of the imaginary part of the free amplitude (Fig. 1) towards the πΣ threshold, which is typical of the single-nucleon approach. Consequently it is reasonable to expect that additional, multi-nucleon terms are required to obtain good fit to the data.

4. \( K^- \)-nucleus optical potentials

Strong interaction level shifts and widths in kaonic atoms have been calculated by solving a Klein-Gordon equation [1] with the optical potential of Eq. (1) transformed to the \( K^- \)-nucleus c.m. system, and where the in-medium \( K^-N \) scattering amplitude is given by the effective amplitude Eq. (8):

\[
V_{K^-} = -\frac{2\pi}{\mu}(1 + \frac{A-1}{A} \frac{\mu}{m_N})F_{K^-N}^{\text{eff}}(\rho)\rho(r),
\]

with \( \mu \) the kaon-nucleus reduced mass and \( \rho = \rho_p + \rho_n \). Two-parameter Fermi distributions were used for both densities, with \( \rho_p \) obtained from the known charge distribution by unfolding the finite size of the charge of the proton. For \( \rho_n \), averages of the ‘skin’ and ‘halo’ forms of Ref. [23] were adopted with the difference between r.m.s. radii given by \( r_n - r_p = (N - Z)/A - 0.035 \text{ fm} \).

Figures 4 and 5 show optical potentials for \( K^-\text{-Ni} \) and for \( K^-\text{-Pb} \), respectively. The potentials marked NLO30 follow directly from the in-medium \( K^-N \) amplitudes in model NLO30, without any adjustable parameters. The agreement with the full data set of 65 points, covering the whole periodic table, is poor, with \( \chi^2 \) per point in the range of 10 – 12. This is not surprising in view of the obvious deficiency of the single-nucleon approach where the imaginary part of the amplitude goes down rapidly towards far subthreshold energies. Adding to the potential an empirical term linear in the nuclear density does not improve much the fit to kaonic atoms data and only by further addition of a \( \rho^2/\rho_0 \) or a \( \rho(\rho/\rho_0)^2 \) term good fits to the data are possible. Both imaginary and real parts of the additional phenomenological potential are then found to be dominated by \( \rho^2 \) or \( \rho^3 \) terms which are likely to represent multi-nucleon absorptive and dispersive contributions, respectively. Figures 4 and 5 also show potentials obtained when adjustable \( b\rho + B\rho^2/\rho_0 \).
terms are added and the four parameters $b$ and $B$ are determined by requiring best fit to the data. The quality of the fits is then quite acceptable, with $\chi^2$ per point of 2 to 2.3. Qualitatively similar results to those displayed in Figs. 4 and 5 are obtained also in the “with SE” version, in agreement with the discussion for Ni in Ref. [12].

Considering values of the potentials at the nuclear center, the additional phenomenological part appears too large to be regarded as a correction term to the basic NLO30 amplitude. However, values of the potential at the center are rather meaningless in the context of kaonic atom observables. The sensitivity of calculated level shifts and widths to the $K^-\,\text{nuclear potentials}$ was found [24] to be around the nuclear surface and certainly not at the center. With that in mind and focusing on the imaginary potential as noted above, it is remarkable that the additional term modifies the shape of the imaginary potentials in the surface region, bringing their r.m.s. radii closer to empirical values. Changes of the imaginary potentials near the surface due to the phenomenological term are of the order of 30% of the NLO30 potentials, consistent with the fraction of multi-nucleon absorption estimated from experiments with emulsions and bubble chambers [25]. The

Figure 4: $K^-\,\text{nuclear potentials}$ for $K^-\,\text{atoms}$ of Ni. Dashed curves: derived from in-medium NLO30 amplitudes; solid curves: plus phenomenological terms from global fits.
emerging phenomenology is similar to that for $V_{\pi^-}$ in pionic atom studies where theoretically motivated single-nucleon contributions are supplemented by phenomenological $\rho^2$ terms representing $\pi NN$ processes [26]. (See also Ref. [1]).

5. Summary and conclusions

A simple ansatz for transforming the strong energy dependence of subthreshold $K^- N$ scattering amplitudes in the nuclear medium to appropriate density dependent averages was presented and employed in global analyses of kaonic atom data, following Refs. [12, 13]. With chirally motivated coupled channel separable interaction scattering amplitudes in model NLO30 [15] that respect the low energy $\bar{K}N$ data, including the recent SIDDHARTA results for kaonic hydrogen [3], the connection between this model and deep real optical potentials was re-established. Effective $K^-$-nucleus amplitudes were derived self-consistently and were critically reviewed with respect to empirical features of phenomenological optical potentials. We focused in the present update on the in-medium effects arising exclusively from the strong energy dependence of subthreshold $K^- N$ amplitudes, using for this purpose
the “without SE” version of the NLO30 in-medium model. The introduction of self-energy effects in the “with SE” version is necessarily model dependent to some extent. Nevertheless, all of our findings and conclusions hold true in both “without SE” and “with SE” versions, with minor differences exhibited already within earlier versions [13]. In the present update, as well as in the preceding studies [12, 13], the steep decrease of the imaginary part of the amplitude as function of the nuclear density, due to the single-nucleon nature of the model, was identified as a major deficiency of the single-nucleon approach. This conclusion is valid also upon adding effective $K^-$-nucleus amplitudes generated by the $p$-wave $\Sigma(1385)$ subthreshold resonance, as discussed in Ref. [13] where $p$-wave effects were found secondary to $\Lambda(1405)$-dominated $s$-wave effects. Good agreement with experiment was achieved by adding to the potential a phenomenological part which was found to be dominated by a $\rho^2$ or a $\rho^3$ term. Including systematically multi-nucleon processes should be the next step in trying to obtain $K^-$-nucleus potentials from in-medium $K^-N$ interaction input.

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