Abstract

We examine the role of the Λ(1405) in kaon-nucleon scattering lengths using chiral perturbation theory. The leading nonanalytic SU(3) corrections reduce the coupling of the Λ(1405) to $K\bar{N}$ compared to $\Sigma\pi$. S-wave $K^-p$ scattering is the only channel significantly affected by the Λ(1405) pole which substantially cancels against the leading term fixed by the vector current. This cancellation and the close proximity of the Λ(1405) pole makes this SU(3) correction to $a(K^-p)$ large and at lowest order leaves the sign of $a(K^-p)$ undetermined. We extract two linear combinations of constants appearing at higher order in the chiral expansion from measurements of $K\bar{N}$ scattering lengths. These constants are important for the offshell behaviour of $K\bar{N}$ scattering amplitudes that play a central role in kaon condensation. The recent claims about the calculability of these constants are discussed.
The effects of the Λ(1405) (denoted below by Λ∗) with \( J^{\pi} = 1/2^- \) on kaon-nucleon scattering lengths have been discussed recently \([1]-[3]\) in relation to the possibility of kaon condensation in dense nuclear matter \([4]-[14]\). The question of whether a kaon condensate forms requires knowledge of the scattering amplitudes of offshell kaons on nucleons at finite density. It is desirable to make a model independent investigation of this phenomenon using only the symmetries of QCD, exploited by the chiral lagrangian. By construction, there will be unknown quantities that must be determined by comparison with experimental data.

We will investigate how the \( KN \) scattering lengths as presently determined constrain the offshell behaviour of these scattering amplitudes. The Λ* spin and parity are such that \( K^- p \) scattering through this resonance occurs in an S-wave and its close proximity to the \( K^- p \) threshold (\( \sim 30 \text{MeV} \)) requires that it be included in any consistent analysis of kaon-nucleon scattering (for the role of the Λ* on kaonic atoms see e.g. \([15]\) and references therein.). Furthermore, its contribution to \( K^- p \) scattering substantially cancels against the leading term arising from the vector current. The pole enhancement of the Λ* contribution and the large cancellation between terms suggests that small changes in the Λ* coupling to \( KN \) may be important. In this letter we compute the leading SU(3) violation in the Λ* coupling to octet baryons in chiral perturbation theory and their effects on kaon-nucleon scattering lengths. Also, at lowest order we extract a linear combination of unknown constants in the offshell scattering amplitudes from the onshell \( K^+ N \) scattering lengths and discuss the recent claims \([2][3][10]\) about calculability of these higher order terms in the chiral expansion.

The dynamics of the pseudo-Goldstone bosons associated with the spontaneous breaking of \( SU(3)_L \otimes SU(3)_R \) chiral symmetry to \( SU(3)_V \) are described by the lagrange density

\[
\mathcal{L}_\pi = \frac{f^2}{8} Tr [\partial_\mu \Sigma^\dagger \partial^\mu \Sigma] + \mu Tr [m_q \Sigma + \text{h.c.}] + \ldots, \tag{1}
\]

where \( \Sigma = \exp(2iM/f) \) is the exponential of the meson field where

\[
M = \begin{pmatrix}
\eta/\sqrt{6} + \pi^0/\sqrt{2} & \pi^+ & K^+ \\
\pi^- & \eta/\sqrt{6} - \pi^0/\sqrt{2} & K^0 \\
K^- & K^0 & -2\eta/\sqrt{6}
\end{pmatrix}, \tag{2}
\]

\( f \) is the meson decay constant (\( f_\pi = 132 \text{MeV} \)) and \( m_q \) is the light quark mass matrix. The dots denote terms higher order in the chiral expansion containing more derivatives or more insertions of the light quark mass matrix. The lowest order chiral lagrangian describing
low momentum interactions of the pseudo-Goldstone bosons with the lowest lying baryons of velocity \( v \) is (for a review see [16])

\[
\mathcal{L}_B = i Tr[\overline{B}_v v \cdot D B_v] - i T^\mu \cdot v \cdot D T_{v \mu} + i \Lambda^*_v v \cdot \partial \Lambda^*_v + \Delta_T T^\mu v T_{v \mu} - \Delta^*_\Lambda^* \Lambda^*_v \Lambda^*_v \\
+ 2 D T_r[\overline{B}_v S^\mu_v \{ A_{\mu}, B \}] + 2 F T_r[\overline{B}_v S^\mu_v \{ A_{\mu}, B \}] + C \left( T^\mu_v A_{\mu} B_v + h.c. \right),
\]

(3)

where \( S^\mu \) is the spin operator, \( A_{\mu} = \frac{i}{2} \left( \xi \partial_{\mu} \xi^\dagger - \xi^\dagger \partial_{\mu} \xi \right) \) is the meson axial current \( (\xi = \exp (iM/f)) \) and \( D B_v = \partial B_v + [V, B_v] \) where \( V_{\mu} = \frac{1}{2} \left( \xi \partial_{\mu} \xi^\dagger + \xi^\dagger \partial_{\mu} \xi \right) \) is the meson vector current. The octet baryons are denoted by the field

\[
B_v = \begin{pmatrix}
\Lambda_v / \sqrt{6} + \Sigma^0_v / \sqrt{2} \\
\Sigma^-_v \\
\Xi^-_v \\
\Xi^0_v \\
-2\Lambda_v / \sqrt{6}
\end{pmatrix},
\]

(4)

and the decuplet of baryon resonances is denoted by the field \( T^\mu_v \). The axial couplings \( F, D, C \) and \( H \) are determined from semileptonic decays of the octet baryons [17] [18] and from the strong decays of the decuplet baryons [19]. The \( \Lambda(1405) \) \( J^\pi = \frac{1}{2}^- \) resonance of four-velocity \( v \) is denoted by \( \Lambda^*_v \) (its inclusion has been discussed by [2]). It has S-wave coupling to a pseudo-Goldstone boson and octet baryon (denoted by \( g_{\Lambda^*} \)) whereas it has only D-wave couplings to a pseudo-Goldstone boson and decuplet baryon, which we will neglect in the work. The mass difference between the octet baryon and the decuplet baryons is denoted by \( \Delta_T \) and the mass difference between the octet baryons and the \( \Lambda^* \) is denoted by \( \Delta^*_\Lambda^* \).

The width of the \( \Lambda^* \) is dominated by the strong decay \( \Lambda^* \to \Sigma \pi \) determined by \( g_{\Lambda^*} \) at lowest order in the chiral expansion. In the limit of exact flavour SU(3) the couplings \( g_{\Lambda^*}, g_{\Lambda^*}(\Sigma \pi) \) (the coupling to \( \Sigma \pi \)) and \( g_{\Lambda^*}(N K) \) (the coupling to \( N K \)) are all equal. However, SU(3) breaking arising from the difference between the mass of the strange quark and the mass of the up and down quarks gives rise to a difference between these couplings. As \( g_{\Lambda^*}(N K) \) is the coupling constant relevant for \( K N \) scattering it is important to estimate the size of this SU(3) breaking. The leading nonanalytic SU(3) violations to the \( \Lambda^* \) couplings arise from the graphs shown in fig. [4]. As the \( \Lambda^* \) is not an asymptotic state we do not include its wavefunction renormalisation in the computation (if we were to treat it as an asymptotic state then it is an SU(3) singlet its wavefunction renormalisation does

\footnote{Our definition of \( g_{\Lambda^*} \) in (3) is a factor of \( \sqrt{2} \) larger than that defined in [2].}
not contribute to SU(3) violating observables at lowest order.) The one-loop coupling constant $g_{\Lambda^*}(BM)$ (coupling of the $\Lambda^*$ to an octet baryon $B$ and pseudo-Goldstone boson $M$) is

$$g_{\Lambda^*}(BM) = g_{\Lambda^*} \left[ 1 + O_M + O_B + O_V(BM) + O_A(BM) \right].$$  \hspace{1cm} (5)

$O_M$ are the contribution from the wavefunction renormalisation of the meson fields fig. 1(a) and are

$$O_n = \frac{2}{96\pi^2 f^2} M_K^2 \log \left( \frac{M_K^2}{\Lambda^2} \right)$$

$$O_K = \frac{5}{96\pi^2 f^2} M_K^2 \log \left( \frac{M_K^2}{\Lambda^2} \right)$$ \hspace{1cm} (6)

while the contribution from the wavefunction renormalisation of the octet baryons fig. 1(b) can be found in [16], and are

$$O_{\Sigma} = -\frac{1}{32\pi^2 f^2} M_K^2 \log \left( \frac{M_K^2}{\Lambda^2} \right) \left[ \frac{26}{3} D^2 + 6F^2 + \frac{14}{3} C^2 \right]$$

$$O_N = -\frac{1}{32\pi^2 f^2} M_K^2 \log \left( \frac{M_K^2}{\Lambda^2} \right) \left[ \frac{17}{3} D^2 + 15F^2 - 10FD + C^2 \right]$$ \hspace{1cm} (7)

We have only retained the nonanalytic contributions from the $K$ and $\eta$ loops using $M_\eta^2 = \frac{4}{3} M_K^2$ and $\Lambda$ is the chiral symmetry breaking scale. For the purposes of this calculation we have neglected the mass splitting between the octet and decuplet baryons $\Delta_T = 0$. The contributions from the graphs involving the two-meson vector coupling arising from the octet baryon kinetic energy term fig. 1(c) are

$$O_V(\Sigma\pi) = \frac{1}{32\pi^2 f^2} \left( M_K^2 F_\beta(\Xi K) + M_K^2 F_\beta(\Lambda K) - 4M_\pi^2 F_\beta(\Sigma\pi) \right)$$

$$O_V(NK) = \frac{1}{32\pi^2 f^2} \left( 3M_K^2 F_\beta(\Xi K) + \frac{3}{2} M_\eta^2 F_\beta(\Lambda\eta) + \frac{3}{2} M_\pi^2 F_\beta(\Sigma\pi) \right)$$ \hspace{1cm} (8)

where the function $F_\beta(BM)$ is shorthand for $F_\beta((M_{\Lambda^*} - M_B)/M_M, M_M/\Lambda)$ and

$$F_\beta(y, z) = (1 - y^2) \log z^2 + 2y \sqrt{y^2 - 1} \log \left( \frac{-y + \sqrt{y^2 - 1 + i\epsilon}}{-y - \sqrt{y^2 - 1 + i\epsilon}} \right)$$ \hspace{1cm} (9)

where we have only retained the leading nonanalytic terms. The last contributions arise from the axial coupling to three mesons fig. 1(d) and are

$$O_A(\Sigma\pi) = -\frac{2}{96\pi^2 f^2} M_K^2 \log \left( \frac{M_K}{\Lambda^2} \right)$$

$$O_A(NK) = -\frac{5}{96\pi^2 f^2} M_K^2 \log \left( \frac{M_K}{\Lambda^2} \right)$$ \hspace{1cm} (10)
which is an equal and opposite contribution to that from $\mathcal{O}_M$ (this is also true for the contributions from the $\pi$’s).

The observed width of the $\Lambda(1405)$ of $\Gamma_{\Lambda^*} = 50 \pm 2$ [20] leads to a value of $|g_{\Lambda^*}(\Sigma \pi)| = 0.58 \pm 0.01$ for $f = f_\pi$. Using $D = 0.7 \pm 0.2$, $F = 0.5 \pm 0.1$, $C = -1.2 \pm 0.2$ and $\Gamma_{\Lambda^*}$ as input parameters for (3) we find that (up to an overall sign)

$$g_{\Lambda^*} = 0.40 \pm 0.04$$

$$g_{\Lambda^*}(\Sigma \pi) = (0.58 \pm 0.01) + (0.12 \pm 0.01)i \quad . \quad (11)$$

$$g_{\Lambda^*}(NK) = (0.32 \pm 0.03) + (0.050 \pm 0.005)i$$

The imaginary parts of the couplings arise from physical intermediate states in the loop graphs shown in fig. 11. The $\Lambda^*$ coupling to $NK$ is seen to be suppressed relative to its coupling to $\Sigma \pi$. It should be remembered that only the leading nonanalytic corrections to $g_{\Lambda^*}(NK)$ and $g_{\Lambda^*}(\Sigma \pi)$ have been included and therefore only part of the SU(3) violation. However, as discussed in [16][19], the leading nonanalytic corrections appear to capture most of the SU(3) violation in axial current matrix elements between octet and/or decuplet baryons and we hope that this is true for these couplings also. The above uncertainties are those associated with the input parameters only. There are systematic uncertainties that are not included arising from truncation of the chiral expansion.

At leading order in the $1/M_N$ expansion ( $M_N$ is the nucleon mass ) and up to order $Q^2$ where $Q = \partial$, $m_s^{1/2} \pm$ the kaon-nucleon S-wave scattering lengths are

$$a(K^\pm p) = \frac{1}{4\pi} \frac{M_N}{M_N + M_K} \left( \mp \frac{2M_K}{f^2} - \frac{g_{\Lambda^*}^2 M_K^2}{f^2} \mp \frac{1}{M_K - \Delta_{\Lambda^*} + i\Gamma_{\Lambda^*}/2} + C_p \right) , \quad (12)$$

$$a(K^\pm n) = \frac{1}{4\pi} \frac{M_N}{M_N + M_K} \left( \mp \frac{M_K}{f^2} + C_n \right) ,$$

where $C_{n,p}$ are unknown constants arising at order $Q^2$ which we will discuss subsequently (at order $Q^3$ and higher the constant $C_p$ for $K^+$ scattering will be different from that for $K^-$ scattering but at order $Q^2$ they are equal). The factor of $M_N/(M_N + M_K)$ is kinematic in origin and is retained explicitly. It comes as no surprise that a new scale, $(M_K - \Delta_{\Lambda^*}) \sim 30\text{MeV}$, has appeared in the $K^-p$ channel from the $\Lambda^*$ resonance, the difference between the position of the pole and the $K^-p$ threshold. Despite the appearance of $M_K^2$ in the contribution from the $\Lambda^*$ suggesting an order $Q^2$ term, the effect of the pole makes it the same size as the leading $M_K$ term in the $K^-p$ channel. Also the finite width
of the $\Lambda^*$ is not higher order for this particular process as its effects are enhanced by this new small scale.

The experimentally determined scattering length for $K^+p$ is $a(K^+p) = -0.31 \pm 0.02\text{fm}$ (for a review of the available data see [21] [22]). The uncertainty in $a(K^+p)$ is an estimate based on fig.2 in [21]. If we set $C_p = 0$ in expression (12) then we would find

$$a(K^+p; C_p = 0, g_{\Lambda^*}) = -0.56 \pm 0.01\text{ fm}$$

which is insensitive to $g_{\Lambda^*}$ and its associated uncertainties as it is far from the $\Lambda^*$ pole ($f = f_\pi$). Neglecting the dependence upon $g_{\Lambda^*}$ we find that

$$\frac{1}{4\pi} \frac{M_N}{M_N + M_K} C_p = 0.25 \pm 0.02\text{ fm}.$$  (14)

Though not surprising, it is somewhat disturbing that this higher order term contributes $\sim 50\%$ of the leading amplitude (it is expected to be smaller than the leading term by only $M_K/\Lambda_\chi \sim 1/2$ from dimensional analysis). It is interesting to ask about predictions we can now make about the $K^-p$ scattering length $a(K^-p)$. Using (12) and the central values for $C_p$ (14) and $g_{\Lambda^*}(BM)$ (11) we find that

$$a(K^-p; g_{\Lambda^*}(\Sigma\pi)) = -0.31 + 0.43i\text{ fm}$$

$$a(K^-p; g_{\Lambda^*}(NK)) = 0.46 + 0.23i\text{ fm}$$

which is very sensitive to the choice of $g_{\Lambda^*}$ due to the close proximity of the $\Lambda^*$ pole and the large cancellation that occurs between this term and the tree-level contribution. It is clear that at this order we are not in a position to make a reliable prediction about either the sign or magnitude of the $K^-p$ scattering length due to the expected large size of corrections. This is in contrast to the results of [2] where the sign of $Re(a(K^-p))$ is found to be negative. Thus $a(K^-p)$ does not provide a good diagnostic with which to test convergence of the chiral expansion and also cannot be used to determine unknown counterterms.

The same procedure can be applied to the kaon-neutron scattering lengths in (12). Scattering amplitudes for $K^+n$ are found from $K^+d$ scattering after removing the proton contribution (see [22]). The extracted scattering lengths for such processes have significant uncertainties and model dependences. For our purposes we assume that $a(K^+n)$ (which has both isospin 0 and 1 contributions) lies between $-0.15$ fm and $-0.30$ fm (estimated...
from fig. 2 in [21]). We fit the constant $C_n$ from this scattering length as there are no pole graphs that can contribute significantly to this amplitude. If we set $C_n = 0$ then we find

$$a(K^+ n ; C_n = 0) = -0.29 \text{ fm}$$

which leads to

$$-0.01 \text{ fm} \lesssim \frac{1}{4\pi} \frac{M_N}{M_N + M_K} C_n \lesssim +0.14 \text{ fm} \ .$$

These in turn lead to predictions for the $K^- n$ scattering lengths at leading order of

$$+0.28 \text{ fm} \lesssim a(K^- n) \lesssim +0.43 \text{ fm}$$

One might be a bit nervous about these predictions for the same reasons that the predictions for the $a(K^- p)$ are unstable to higher order corrections. However, the nearest confirmed S-wave $I = 1$ resonance is the $\Sigma(1620)$ ($J^\pi = \frac{1}{2}^-$) which only weakly couples to $NK$ [20]. Therefore we may expect that the meson interactions calculable in chiral perturbation theory are the leading corrections to our results.

If we are interested in possibility of kaon condensation in dense nuclear matter then it is the scattering amplitudes for offshell kaons that are required. In order to make progress in this direction we need to understand the origin of the constants $C_{n,p}$ appearing in (12). At order $Q^2$ the lagrange density responsible for $C_{n,p}$ is given by

$$\mathcal{L}_{\text{mass}} = a_1 \text{Tr}[\overline{B}_v \chi_+ B_v] + a_2 \text{Tr}[\overline{B}_v B_v \chi_+] + a_3 \text{Tr}[\overline{B}_v B_v] \text{Tr}[\chi_+]$$

$$+ d_1 \text{Tr}[\overline{B}_v A^2 B_v] + d_2 \text{Tr}[\overline{B}_v (v \cdot A)^2 B_v] + d_3 \text{Tr}[\overline{B}_v B_v A^2]$$

$$+ d_4 \text{Tr}[\overline{B}_v B_v (v \cdot A)^2] + d_5 \text{Tr}[\overline{B}_v B_v] \text{Tr}[A^2] + d_6 \text{Tr}[\overline{B}_v B_v] \text{Tr}[(v \cdot A)^2]$$

$$+ d_7 \text{Tr}[\overline{B}_v A^\mu] \text{Tr}[A_\mu B_v] + d_8 \text{Tr}[\overline{B}_v (v \cdot A)] \text{Tr}[(v \cdot A) B_v]$$

where $a_{1,2,3}$ and $d_{1...8}$ are unknown constants that must be determined experimentally and $\chi_+ = \xi m_q \xi + \xi^\dagger m_q \xi^\dagger$. The coefficients $a_1$ and $a_2$ can be determined from the octet baryon mass splittings [23] where it is found that $m_s a_1 = -65\text{MeV}$ and $m_s a_2 = -125\text{MeV}$. $a_3$ is the “sigma term” which does not contribute to mass splittings but can be measured by low energy $\pi N$ scattering where it is found to be sensitive to kaonic loop corrections [23] [24]. It is important to stress that the coefficients $d_i$ are not calculable from chiral symmetry arguments alone and in particular there is no sense in which they can be computed as has
been claimed in [2] and discussed in [10][3].

The constants $C_{p,n}$ determined by (14) constrains two linear combination of the $a_i$ and $d_i$ given by

$$C_p = -\frac{2m_s}{f^2} (a_1 + a_2 + 2a_3) + \frac{M_K^2}{f^2} (d_1 + d_2 + d_3 + d_4 + 2d_5 + 2d_6 + d_7 + d_8),$$

$$C_n = -\frac{2m_s}{f^2} (a_2 + 2a_3) + \frac{M_K^2}{f^2} (d_3 + d_4 + 2d_5 + 2d_6)$$

where $m_s$ is the mass of the strange quark and we have neglected the mass of both the up and down quarks. The constants $d_i$ have dimensions of inverse mass, and naive dimensional analysis would indicate $d_i \sim 1/\Lambda_\chi$. At the order to which we are working these are the only two constraints that can be obtained from onshell $KN$ scattering. The offshell $KN$ scattering amplitudes become a function of $a_3$ alone (neglecting loop and finite density effects) since the same linear combinations of $d_i$ and $a_i$ appear both onshell and offshell and $a_{1,2}$ are fixed from octet baryon mass splittings. Low energy pion-nucleon scattering is described by the lagrange densities given in (3)(19) and a recent analysis of $\pi N$ scattering lengths can be found in [25]. The constants $d_{3,4,7,8}$ and $a_2$ do not contribute to these scattering amplitudes and hence cannot be constrained by the experimentally determined $\pi N$ scattering lengths.

In conclusion, we have computed the leading nonanalytic SU(3) violating corrections to the $\Lambda^*$ coupling constant. We find that the $\Lambda^*$ coupling to $KN$ is suppressed relative to its coupling to $\Sigma\pi$. The substantial cancellation between the tree-level and $\Lambda^*$ contributions to the $K^-p$ scattering amplitude and the close proximity of the $\Lambda^*$ pole to the $NK$ threshold enhances the effect of this SU(3) correction. We fit two linear combinations of coefficients of higher dimension operators to $a(K^+p)$, $a(K^+n)$ and find that the sign of $a(K^-p)$ cannot be determined at this order contrary to the result of [2]. We also stress that the coefficients of higher dimension operators required to compute offshell scattering amplitudes important for kaon condensation cannot be calculated from chiral symmetry arguments. Corrections to our results arising from loop graphs and higher dimension operators are expected to be suppressed by only $M_K/\Lambda_\chi \log (M_K^2/\Lambda_\chi^2)$ compared to the $C_{n,p}$ and may be significant.

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2 In [2][3][11] the contribution to the coefficients $d_i$ arising at order $1/M_N$ are computed from the “relativistic” lagrangian for nucleon interactions. It is important to realise that there are operators, e.g. $Tr[D_\mu B_\nu A^\mu A^\nu D_\nu B_\nu]$ that should be included in addition to the contributions arising at order $1/M_N$. Such operators have unknown coefficients and give a contribution of order $(M_N/\Lambda_\chi)^2$ for nucleons near their mass shell, in fact operators with arbitrary powers of $D^\mu$ will give unsuppressed contributions.
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Figure Captions

Fig. 1. The Feynman diagrams giving rise to the leading SU(3) violation in the coupling of the Λ(1405) to octet baryons and pseudo-Goldstone bosons. The dashed lines denote pseudo-Goldstone bosons and the solid lines denote baryons. Graphs (a) and (b) are wavefunction renormalisations for the mesons and baryons respectively and graphs (c) and (d) are vertex corrections.
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