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Sifting through the SM for the hints of an ALP

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We show how a light axion-like particle (ALP) leaves tell-tale signatures in processes that involve standard model (SM) fields only (i.e., SM processes). These include the violation of the Gell-Mann–Okubo mass relation, modification of form factors, alteration of differential rates for various SM transitions, gives rise to novel sum rules etc. Therefore, in the presence of an ALP, extractions of masses, mixing angles, and form factors in a data-driven way provide important (indirect) bounds on ALP physics. These bounds remain valid in the limits where new physics effects conspire to weaken the bounds from direct searches. We provide a proof of concept example by analysing $K^+_{\ell_3}$ decays. We also derive sum rules which give hints towards the nature of the ALP physics in the UV.

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1. Introduction

Following the work discussed in Ref.[1], we will discuss some examples of how the chiral Lagrangian of the Standard Model (SM) is modified in the presence of an Axion-like particle (ALP, a) and then go on to talk about how we can use these modifications to give indirect constraints on the Wilson coefficients of the effective Lagrangian.

The most straightforward way of finding new physics (NP) is to look for the presence of new particles in the final states of decays or collisions. The hope is to find spectacular signatures like bumps, missing E_T^{miss}, displaced tracks, disappearing tracks etc. Although straightforward, one issue in using these channels is that these methods depend crucially on the physical properties of the particle being looked for, e.g., mass, lifetime, branching ratios etc. Different experiments are sensitive to different regions of the parameter space, with some regions being beyond the scope of current or upcoming experiments. It is here that indirect channels come to be of use. Indirect channels only have SM particles in the final state, and are not highly sensitive to the physical properties of the NP particle.

Now, there is a flip side. When working with direct detection techniques, one can just write one allowed operator, say $aF\widetilde{F}$, and give bounds on the corresponding Wilson coefficient. However, when looking for modifications to SM processes in the presence of NP, we have to write down all operators as allowed by symmetry and check the contributions from all the different Wilson coefficients. So, it becomes an exercise of writing an EFT Lagrangian, with the SM and the NP degrees of freedom. This is what we do first.

2. Formalism

In order to derive the chiral Lagrangian in the presence of the ALP, we first write the quark–ALP Lagrangian in a convenient form:

$$\mathcal{L} \supset \overline{q}_L\gamma^\mu\left(i\partial_\mu + L_\mu\right)q_L + \overline{q}_R\gamma^\mu\left(i\partial_\mu + R_\mu\right)q_R + \overline{q}_L\overline{M}q_R + \cdots \text{ , where}$$
$$L^\mu = L^\mu_{\text{SM}} + \frac{\partial^\mu a}{f_a}C_L^8 t_8\,,\ R^\mu = R^\mu_{\text{SM}} + \frac{\partial_\mu a}{f_a}C_R^i t_i\,,\ \overline{M} = \left(1 + iC^i_{LR}t_i\frac{a}{f_a} + \cdots\right)M. \tag{1}$$

Here, $q_{L/R} \equiv (u, d, s)_{L,R}$. Repeated indices are summed over and t_i are the generators of $SU(3)$ of flavor. We nominally gauge $SU(3)_L \times SU(3)_R$ of flavor, with L_μ, R_μ the corresponding compensating gauge transformations of the quark fields. Treating the mass term $\overline{M} \to L\overline{M}R^\dagger$ as a spurion, even it acts as a 'gauge-invariant' term. For brevity, we don't give the well-known forms of M (quark mass matrix), L^μ_{SM}, and R^μ_{SM}. We do not add electroweak (EW) breaking operators, also, we avoid writing operators giving rise to flavor changing neutral currents (FCNC).

With the quark–ALP Lagrangian written in this form, it is straightforward to derive the modified chiral Lagrangian. In terms of the exponential representation of the pions, $U_\pi \equiv \exp(2i\pi^a t_a/f_\pi)$, we have, after current matching:

$$\mathcal{L} \supset \frac{f_\pi^2}{4}\,\text{Tr}\Big[\left|\partial_\mu U_\pi - i(L_\mu U_\pi - U_\pi R_\mu)\right|^2\Big] + \frac{\Lambda f_\pi^2}{2}\,\text{Tr}\Big[\overline{M}U_\pi^\dagger\Big] + \text{h.c.} + \cdots \tag{2}$$

Here, Λ is the UV cut off of the (chiral-)EFT, and the $\cdots$ represent higher order terms in the chiral Lagrangian. The Lagrangian in eq. (2), with L_μ, R_μ, and M given by eq. (1), gives the leading

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order terms in the chiral Lagrangian with the ALP (AχPT). The presence of the ultraviolet (UV) suppression scale of the ALP, f_a, introduces a new power counting parameter into the AχPT. We use $\xi \equiv f_\pi/f_a$ as the power counting parameter in this work and work till $O(\xi^2)$.

All our results are obtained by expanding the Lagrangian in eq. (2). We will discuss three of the most notable ones here, the first of these is the modification to the meson mass spectrum. The ALP has both kinetic and mass mixing terms with the π^0 and the η. Therefore, the physical π^0 and the physical η are redefined in the presence of the ALP. The modified masses in the AχPT are:

$$M^2_{\pi^0} = 2B_0\hat{m}\left[1+\frac{\xi^2}{6}\left(3C_3^2 - 2\sqrt{3}C^8_{LR}C^3_R\frac{m_\Delta}{\hat{m}}\right)\right], \qquad M^2_\eta = \frac{4}{3}B_0\left(m_s+\frac{1}{2}\hat{m}\right)\left[1+\frac{\xi^2}{4}C_8^2\right]. \quad (3)$$

The other masses are the same as in the SM, and $B_0 \equiv \Lambda, m_\Delta \equiv (m_u - m_d)/2, \hat{m} \equiv (m_u + m_d)/2$, $C_3 \equiv C^3_{LR} - C^3_R, C_8 \equiv C^8_{LR} - C^8_A$. Using the deviation of the trajectories of the meson masses, we obtain these modifications (at leading order (LO)) in the following sums of the meson masses:

$$\Delta_{M_\pi} \equiv \frac{M^2_{\pi^+} - M^2_{\pi^0} - \Delta_e}{M^2_{\pi^+}} = 0 + \frac{1}{6}\left(2\sqrt{3}C^8_{LR}C^3_R\frac{m_\Delta}{\hat{m}} - 3C_3^2\right)\xi^2 + \cdots\ , \quad (4a)$$

$$\Delta_{\mathrm{GMO}} \equiv \frac{4M_K^2 - M_\pi^2 - 3M_\eta^2}{M_\eta^2 - M_\pi^2} = 0 - \frac{3}{4}\xi^2 C_8^2 + \cdots\ . \quad (4b)$$

Here, Δ_e is the electromagnetic (EM) correction to the charged pion mass. We use the isospin invariant definitions $M_K^2 = \left(M^2_{K^0} + M^2_{K^\pm}\right)/2$, $M_\pi^2 = M^2_{\pi^\pm}$. Both these sums are zero in the SM (up to $\eta - \eta'$ mixing). However, in the AχPT, the sum relations get violated at the tree level itself.

Operator, $O^i_{K^+_{\ell_3}}$	Coefficient, $C^i_{K^+_{\ell_3}}$
$O^0_{K^+_{\ell_3}} = [K^+\partial_\mu(\pi_0+\sqrt{3}\eta) - \partial_\mu K^+(\pi_0+\sqrt{3}\eta)]j^\mu_{-,\ell}$	$C^0_{K^+_{\ell_3}} = iG_F V_{\bar{s}u}$
$O^1_{K^+_{\ell_3}} = \left(K^+\partial_\mu a - \partial_\mu K^+ a\right) j^\mu_{-,\ell}$	$C^1_{K^+_{\ell_3}} = iG_F V_{\bar{s}u}\frac{\xi}{2}(C^3_R + \sqrt{3}C^8_R)$
$O^2_{K^+_{\ell_3}} = \left(K^+\partial_\mu a + \partial_\mu K^+ a\right) j^\mu_{-,\ell}$	$C^2_{K^+_{\ell_3}} = iG_F V_{\bar{s}u}\frac{\xi}{2}(C^3_R + \sqrt{3}C^8_R)$
$O^3_{K^+_{\ell_3}} = \partial^\mu a\ \left(\partial_\mu K^+K^- - K^+\partial_\mu K^-\right)$	$C^3_{K^+_{\ell_3}} = \frac{i}{4}\frac{1}{f_\pi}\xi\left(C^3_R + \sqrt{3}C^8_V\right)$
$O^4_{K^+_{\ell_3}} = \partial_\mu K^+ j^\mu_-$	$C^4_{K^+_{\ell_3}} = -2f_\pi\, G_F V_{\bar{s}u}$

Table 1: Operators in the *original-basis* (before any mixing) contributing to $K^+_{\ell_3}$ decay at tree level.

The second modification is to the meson decay amplitudes, of which we discuss the semi-leptonic charged current (CC) decays in this work. The $\pi^+ \to \pi^0\ell^+\nu$ (π_β) and the $K^+ \to \pi^0\ell^+\nu$ ($K^+_{\ell_3}$) decays are driven exclusively by the $(\partial_\mu K^+(\pi^+)\pi^0 - K^+(\pi^+)\partial_\mu\pi^0)\ell\gamma^\mu\bar{P}_L\nu$ operators at leading order in SMχpt. However, in the AχPT, there are multiple terms that are relevant for these decays, even at leading order, as listed in table 1 (with $j^\mu_{-,\ell} = \bar{\nu}\gamma^\mu\frac{1}{2}(1-\gamma_5)\ell$).

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The operators in table 1 are in the flavor basis. The *eigenbasis* is obtained by rotating the fields to get rid of kinetic- and mass-mixing. These new operators modify the form factors (FF) governing the amplitude of $K^\pm \to \pi^0 \ell^\pm \nu$,

$$\begin{aligned}
\mathcal{A} &= G_F V_{\bar{s}u} \left[\tilde{f}_+^{K^+\pi^0}(0)\, Q_\mu + \tilde{f}_-^{K^+\pi^0}(0)\, q_\mu \right] \bar{u}_\nu \gamma^\mu \frac{1}{2} (1-\gamma_5)\, v_\ell. \quad \text{where,} \\
\tilde{f}_+^{K^+\pi^0}(0) &= \alpha^{(0)}_{K^+\pi_0} + \xi^2 \alpha^{(2)}_{K^+\pi_0}; \quad \tilde{f}_-^{K^+\pi^0}(0) = \beta^{(0)}_{K^+\pi_0} + \xi^2 \left(\beta^{(2)}_{K^+\pi_0} + i\tilde{\beta}^{(2)}_{K^+\pi_0} \right),
\end{aligned} \tag{5}$$

$$\alpha^{(0)}_{K^+\pi_0} = 1 - \sqrt{3}\,\epsilon,\ \alpha^{(2)}_{K^+\pi_0} = -\frac{C_3}{8}(C^3_{LR} - C^3_R + 2\sqrt{3}(C^8_{LR} - C^8_R)),\ \beta^{(0)}_{K^+\pi_0} = 0,\ \beta^{(2)}_{K^+\pi_0} = -\frac{\sqrt{3}}{4} C_3 C^8_L\,.$$

There are additional contributions coming from higher orders in the chiral expansion, from electromagnetism, and from EW breaking operators [2–4]. We absorb the higher order corrections coming from EW, EM effects, and higher order terms in the chiral expansion into $\tilde{f}_\pm^{K^+\pi^0}(0) \to \tilde{f}_\pm^{K^+\pi^0}(t)$ where $t \equiv q^2$. With, all these in hand, we can finally express the amplitude-squared as:

$$\begin{aligned}
\overline{|\mathcal{A}|}^2_{K_{l3}} &= 2G_F^2 |V_{\bar{s}u}|^2 C_{\text{cor}} \left[1 + 2\,\xi^2 \frac{\alpha^{(2)}_{K^+\pi^0}}{\alpha^{(0)}_{K^+\pi^0}} \right] (2H\cdot p_\ell\ H\cdot p_{\nu_\ell} - H^2 p_\ell \cdot p_{\nu_\ell}), \\
\text{where } H_\mu &\equiv f^{K^+\pi^0}_{+,\text{SM}}(t)\, Q_\mu + \left[1 + \xi^2 \left(\frac{\beta^{(2)}_{K^+\pi^0}}{\delta\beta^{(0)}_{K^+\pi^0}} - \frac{\alpha^{(2)}_{K^+\pi^0}}{\alpha^{(0)}_{K^+\pi^0}} \right) \right] f^{K^+\pi^0}_{-,\text{SM}}(t)\, q_\mu.
\end{aligned} \tag{6}$$

The factor C_{cor} represents effects which are not included in the lattice computations of the FFs that we use. Next, we use $K^\pm \to \pi^0 \ell^\pm \nu$ decay rates and distributions to bound $\xi^2 \alpha^{(2)}_{K^+\pi^0}$ and $\xi^2 \beta^{(2)}_{K^+\pi^0}$.

3. Constraints from K_{ℓ_3} decay

To constrain $\xi^2 \alpha^{(2)}_{K^+\pi^0}$ and $\xi^2 \beta^{(2)}_{K^+\pi^0}$, we use the following independent measurements:

- Measurement of the differential decay distributions of $K^\pm \to \pi^0 \mu^\pm \nu_\mu (K^+_{\mu_3})$ and $K^\pm \to \pi^0 e^\pm \nu_e (K^+_{e_3})$ by the NA48/2 collaboration at the CERN SPS [5].
- The total width measurements for $K^+_{\mu_3}$ and $K^+_{e_3}$ decays. We have used the experimental averages of the branching fractions from PDG [6] to calculate the rates.

As is evident from eq. (6), the effect of $\xi^2 \alpha^{(2)}_{K^+\pi^0}$ dominates over that of the lepton-mass–suppressed effect of $\xi^2 \beta^{(2)}_{K^+\pi^0}$ when it comes to the total decay rate. However, as we show below, the marginal energy spectra of the decay rate of $K^+_{\mu_3}$ can be used effectively to constrain $\xi^2 \beta^{(2)}_{K^+\pi^0}$.

For the SM contributions to the FF parameters, we use results obtained from lattice computations by the European twisted mass collaboration [7]. On the experiment side, we use the $K^+_{\ell_3}$ data obtained by the NA48/2 collaboration to fit the truth level distribution against observation. The data consist of bin-by-bin event distributions of the differential decay rate with respect to the pion and the lepton energies (E_π, E_μ), for 4.4×10^6 and 2.3×10^6 reconstructed events corresponding to $K^+_{e_3}$ and $K^+_{\mu_3}$ respectively. We combine this multi-variable fit with the constraints set by the independent measurements of the total decay rates to bound $\xi^2 \alpha^{(2)}_{K^+\pi^0}$ and $\xi^2 \beta^{(2)}_{K^+\pi^0}$.

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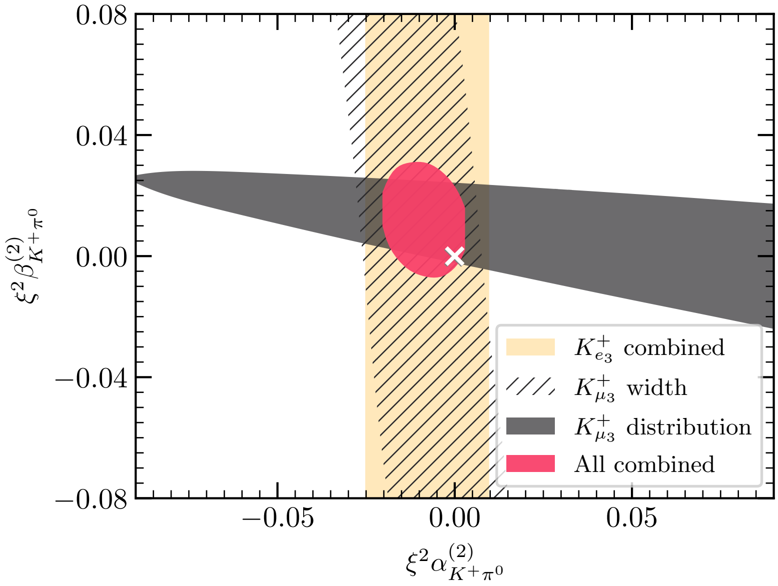

Figure 1: The 95% C.L. allowed regions for the ALP parameters in the $\xi^2\alpha^{(2)}_{K^+\pi^0}$ - $\xi^2\beta^{(2)}_{K^+\pi^0}$ plane. The yellow band indicates the region allowed by the combined $K^+_{e_3}$ data, *i.e.*, the total rate and differential rates combined. The black patch shows the region allowed (at 95% C.L.) by the $K^+_{\mu_3}$ differential distribution. The hatched area is the corresponding region allowed by the $K^+_{\mu_3}$ total rate. The red patch is the 95% C.L. allowed region obtained by combining all the independent analyses. The cross marks the SMχPT point where the values of both the parameters are zero.

We compute χ^2 distributions by comparing the truth-level signal against the differential distribution data and the total decay width measurement, after taking into account the experimental and theoretical errors and correlations. For the differential distributions, we normalize our histograms using the total number of events, as quoted in the last paragraph. In Figure 1, we show the 95% confidence limits (C.L.) obtained in our analysis in the ($\xi^2\alpha^{(2)}_{K^+\pi^0}$ - $\xi^2\beta^{(2)}_{K^+\pi^0}$) plane.

4. Sum Rules in meson decays

Using the modified form factors defined above, we can construct sum rules (in the $m_\ell \to 0$ limit) which can show deviations from the SM expectations. In the standard model, due to the completeness of the π^0–η basis, we have:

$$\frac{1}{4}\left|f^{K^+\pi^0}_{+,\,\mathrm{SM}}(0)\right|^2 + \frac{3}{4}\left|f^{K^+\eta}_{+,\,\mathrm{SM}}(0)\right|^2 = 1. \tag{7}$$

When the complete basis involved the ALP, it is reflected in this sum as well. In the AχPT, in terms of the modified FFs, we have:

$$\frac{1}{4}\left|\tilde{f}^{K^+\pi^0}_{+}(0)\right|^2 + \frac{3}{4}\left|\tilde{f}^{K^+\eta}_{+}(0)\right|^2 = 1 - \frac{\xi^2}{16}\left(C^3_{LR} - C^3_R + \sqrt{3}(C^8_{LR} - C^8_R)\right)^2 + \xi^2\frac{3}{16}(C^8_L)^2 \tag{8}$$

For $C^8_L = 0$, the sum is identically less than one, as expected from considerations of completeness. However, when there are t^8_L breaking interaction between the ALP and the mesons, i.e. for $C^8_L \neq 0$,

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this sum can be greater than one. Therefore, a positive deviation of the sum from unity is possible and it uniquely signals a t_L^8 breaking interaction between the ALP and the quarks in the UV. Therefore, this sum can tell us about the corresponding UV model. The latter would not have been possible by just looking at the deviations from SM expectations of the individual decay widths .

5. Conclusions

We have discussed a few of the avenues in which the presence of a low-lying ALP can be identified by searching for deviations from the SM expectations in observables with mesons and leptons in the final states. For brevity, only a few findings have been discussed and more such results can be found in Ref. [1]. One particularly interesting result not discussed here is the orthogonality of the bounds obtained from these indirect method to direct detection bounds. As an example, it should be mentioned that the functional form (in terms of Wilson coefficients) of the deviation of the $K^\pm \to \pi^0 \ell^\pm \nu$ amplitude from the SM expectation is different from the functional form of the $K^\pm \to a \ell^\pm \nu$ amplitude. Therefore, the limit in which one goes to zero is not the same as that for the other going to zero. Hence, this is a way to look for ALPs for the so called 'pion-phobic' models [8]. As there are much more flavor observable than what we discuss here, we can expect the formalism discussed in this work to be used in a more general context, for more flavor variables to get novel bounds on the ALP parameter space.

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