Probing the Light Pseudoscalar Window

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(Dated: March 2003)

Abstract

Very light pseudoscalars can arise from the symmetry-breaking sector in many extensions of the Standard Model. If their mass is below 200 MeV, they can be long-lived and have interesting phenomenology. We discuss the experimental constraints on several models with light pseudoscalars, including one in which the pseudoscalar is naturally fermiphobic. Taking into account the stringent bounds from rare $K$ and $B$ decays, we find allowed parameter space in each model that may be accessible in direct production experiments. In particular, we study the photoproduction of light pseudoscalars at Jefferson Lab and conclude that a beam dump experiment could explore some of the allowed parameter space of these models.
I. INTRODUCTION

In many extensions of the standard model, the electroweak symmetry-breaking sector includes additional weak doublets or singlets. New CP-even, CP-odd and charged scalar states may be present in the physical spectrum. The masses of these particles are typically of the same order as the weak scale, and fine-tuning is required to make them much lighter. An exception occurs if the theory possesses an approximate global symmetry: a CP-odd scalar may become a massless goldstone boson in the limit that such a symmetry is exact, and a massive state that is naturally light in the case where the symmetry is only approximate. We will henceforth refer to such CP-odd states as light pseudoscalars.

The most familiar example of a light pseudoscalar is the axion. This pseudo-goldstone boson arises in a two-Higgs-doublet model with a global symmetry that allows independent phase rotations of the two Higgs fields. The axion arises as a consequence of spontaneous symmetry breaking and is exactly massless in the absence of gauge interactions. The axion acquires a small mass due to the QCD anomaly, which breaks this global symmetry at the quantum level.

In other models, a global symmetry may be broken more significantly by a small parameter that appears explicitly in the Lagrangian. For example, consider the Higgs potential for two Higgs doublets, with a \( \Phi_2 \leftrightarrow -\Phi_2 \) symmetry:

\[
V = \mu_1^2 \Phi_1^\dagger \Phi_1 + \mu_2^2 \Phi_2^\dagger \Phi_2 + \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \lambda_4 |\Phi_1^\dagger \Phi_2|^2 + \frac{\lambda_5}{2} ((\Phi_1^\dagger \Phi_2)^2 + \text{h.c.})
\]

(1.1)

In the limit \( \lambda_5 \to 0 \), this potential has a \( U(1) \times U(1) \) symmetry in which each doublet rotates by an independent phase. The spontaneous breaking of the diagonal \( U(1) \) symmetry yields a goldstone boson that is “eaten” when the theory is gauged; the remaining \( U(1) \), which rotates each doublet by an opposite phase, yields a physical goldstone boson state. When \( \lambda_5 \) is nonvanishing, this pseudoscalar develops a mass given by \( m_A^2 = -\lambda_5 v^2 \), where \( v = 246 \text{ GeV} \) is the electroweak scale. In this paper, we will consider pseudoscalars with masses in the \( 100 - 200 \text{ MeV} \) range, for phenomenological reasons explained below. This can be achieved by setting \( \lambda_5 \) equal to a small number that is comparable to a light fermion Yukawa coupling—a light pseudoscalar would then be no more or less unnatural than a muon or light quark.
Of course, one can construct models in which the light fermion Yukawa couplings arise only via higher-dimensions operators in a more complete high-energy theory. The Yukawa couplings are identified with powers of the ratio of a symmetry breaking scale to the cut off of the theory, and therefore can be naturally small. By analogy, the U(1) symmetry present in the $\lambda_5 = 0$ limit of Eq. (1.1) may be broken by a field $\eta$ that acquires a vacuum expectation value at some high scale and contributes to the term of interest only through Planck-suppressed operators. Given this dynamical assumption, one predicts that the pseudoscalar mass is of the order $(\langle \eta \rangle / M_*)^{n/2} v$, where $n$ is a positive integer, and $M_* = 2 \times 10^{18}$ GeV is the reduced Planck mass. Interestingly, for $n = 2$, and $\langle \eta \rangle \sim 10^{15}$ GeV (the nonsupersymmetric GUT scale), one obtains a pseudoscalar mass of approximately 100 MeV. One can imagine a variety of high energy theories in which similar results are obtained.

Our interest in pseudoscalar masses between 100 and 200 MeV is motivated by the pseudoscalar decay length and production cross section. We hope to have both in optimal ranges for detection of the pseudoscalar in possible photoproduction experiments at Jefferson Lab. As far as production is concerned, existing direct searches yield bounds on the pseudoscalar couplings that are weakest in this mass range, and a wide variety of experiments severely constrain the pseudoscalar couplings for masses below 100 MeV. On the other hand, if the pseudoscalars are produced in significant but not overwhelming numbers, we hope for a decay length that is long enough to clearly separate the pseudoscalar decay signal from possible mesonic backgrounds. Pseudoscalars with masses above 200 MeV decay rapidly into muon pairs with a branching fraction near 100%, making detection via a separated vertex impossible. Thus, the $100 – 200$ MeV mass window seems particularly promising for the experimental search that we propose in Section IV.

To proceed with our phenomenological analysis, we must decide on the pseudoscalar’s couplings to standard model fermions; the pattern of these couplings is in fact quite model-dependent. In the standard two-Higgs-doublet models, the pseudoscalar couplings are proportional to Yukawa matrices multiplied by a ratio of the vacuum expectation values $v_1$ and $v_2$. On the other hand, one can employ simple discrete symmetries to construct three-doublet models in which only two doublets couple to quarks and do not mix with a third doublet coupling to the leptons. In this case, the pseudoscalar in the quark-two-doublet sector is entirely leptophobic. An analogous three-doublet model with a lepton-two-doublet sector
yields a pseudoscalar that has no couplings to quarks and is, hence, hadrophobic. Such models illustrate the range of the possible, but are not particularly well motivated. A much more appealing possibility is that the pseudoscalar may have no direct couplings to quarks or leptons at all. Let us comment on the motivation for such a fermiophobic pseudoscalar in more detail.

One could imagine a number of reasons why a pseudoscalar may have suppressed couplings to standard model fermions. The suppression could be parametric, as in the type-I two Higgs doublet model when \( \tan \beta \) is taken large. On the other hand, the suppression could be geometric, as in extra-dimensional scenarios in which fields have wave functions that are localized at different points in an extra dimension. Let us focus on a concrete realization of this second idea. Consider an \( S^1/Z_2 \) orbifold of radius \( R \), with standard model matter fields located at the \( y = 0 \) fixed point, and gauge fields in the 5D bulk. Here \( y \) is the extra-dimensional coordinate. Assume that there exists additional vector-like matter in complete SU(5) representations (to preserve gauge coupling unification) as well as a gauge-singlet scalar field \( S \), all isolated at the \( y = \pi R \) fixed point. A spontaneously broken approximate global symmetry of the singlet potential leads to a light pseudoscalar state that couples directly to the exotic matter multiplets only. The geometry of this scenario prevents mixing between the ordinary and exotic matter fields, which communicate with each other only via gauge interactions in the bulk. The scale of compactification can be taken large enough so that the effects of Kaluza-Klein excitations are irrelevant to the low-energy theory.

Given the simplicity of the fermiophobic singlet scenario described above, we will focus our discussion on light pseudoscalars in the two-Higgs-doublet models of type-I and II and in the fermiophobic singlet scenario. We comment on the other possibilities where appropriate. In Section II, we analyze the experimental constraints on the light pseudoscalar in the conventional two-Higgs-doublet models, placing particular emphasis on the bounds from \( K \) and \( B \) meson decays. In Section III, the fermiophobic singlet scenario is studied, and in Section IV we study the possibility of detecting pseudoscalars of either type in photoproduction experiments at Jefferson Lab. Section V contains our conclusions.
II. CONSTRAINTS IN TWO-DOUblLET MODELS

As we have described in the previous section, light pseudoscalars can arise in two-Higgs-doublet extensions of the standard model. Two popular options exist in which a discrete symmetry is imposed to forbid tree-level flavor changing neutral currents [7]: In Model I, all of the fermions couple to a single Higgs doublet, but none to a second. In Model II, the charge \( Q = 2/3 \) quarks couple to one Higgs doublet while the \( Q = -1/3 \) quarks and the leptons couple to another. A third possibility is that all fermions couple to both Higgs doublets, without the restriction of any discrete symmetry. An ansatz is then employed to make tree-level flavor changing Higgs couplings sufficiently small [8]. However, in this case it has been shown that a very light pseudoscalar will still lead to unacceptably large flavor-changing neutral currents [9].

The coupling of the pseudoscalar Higgs to fermions is of the form \(- \frac{m_f}{v} X_f \bar{f} \gamma_5 f A\) where \(v = 246\) GeV and \(X_f = \cot \beta\) for all fermions in Model I, and \(X_f = \cot \beta \left(\tan \beta\right)\) for the \(Q = 2/3\) quarks \((Q = -1/3\) quarks and leptons) in Model II. Here \(\tan \beta\) is the ratio of vacuum expectation values of the two Higgs doublets, and is a free parameter.

There have been numerous discussions of the bounds on a light pseudoscalar, most recently by Larios, Tavares-Velasco and Yuan [10]. In Model II, the combined bounds from the nonobservation of \(J/\Psi \rightarrow A\gamma\) and \(\Upsilon \rightarrow A\gamma\) force \(\tan \beta\) to be close to 1, since the former decay implies \(\tan \beta \lesssim 1\) and the latter implies \(\cot \beta \lesssim 1\) [3]; theoretical uncertainties don’t quite allow the model to be excluded. In Model I, both decays imply only that \(\cot \beta \lesssim 1\).

Bounds from \(\eta, \eta'\) and \(\pi\) decays also force \(\tan \beta \sim 1\) in Model II and \(\cot \beta \lesssim 1\) in Model I [11]. Bounds from \(g - 2\) are in flux at the moment, but do not appreciably change these results. (In addition, the \(g - 2\) bound is only valid if one makes a strong assumption that there are no other possible nonstandard contributions at one loop.) Bounds from \(b \rightarrow s\gamma\), \(\Delta \rho\), \(R_b\) and \(A_b\) can all be avoided by constraining the neutral and charged scalar masses [10]. Thus, we will consider two cases: Model II with \(\tan \beta \sim 1\) and Model I with \(\cot \beta \lesssim 1\). After reviewing the decay modes and decay lengths of the light pseudoscalar, we consider the bounds from \(K\) and \(B\) meson decays, which present the strongest constraints on these models.
A. Decay Modes

For a pseudoscalar lighter than twice the muon mass, there are only two possible decay modes, $A \rightarrow e^+e^-$ and $A \rightarrow \gamma\gamma$. The decay width into an electron pair is given by

$$\Gamma_{A \rightarrow e^+e^-} = \frac{m_e^2}{8\pi v^2} M_A X_e^2 \left(1 - 4 \frac{m_e^2}{M_A^2}\right)^{1/2}. \tag{2.1}$$

For $\tan\beta = 1$, this gives a decay length of $0.6 - 1.2$ centimeters in the pseudoscalar rest frame, for $M_A$ ranging from 100 to 200 MeV. This result scales as $\tan^2\beta$ in Model II and $\cot^2\beta$ in Model I.

The decay into two photons proceeds at one loop with the width

$$\Gamma_{A \rightarrow \gamma\gamma} = \frac{\left|\sum_f N_c Q_f^2 X_f\right|^2 \alpha^2 M_A^3}{64\pi^3 v^2}, \tag{2.2}$$

where $N_c$ is 3 for quarks and 1 for leptons and $Q_f$ is the fermion charge. This expression is valid if the mass of the fermion in the loop is much larger than the momentum in the decay. When this is not the case then the exact expression given in Refs. [3, 10] should be used. Note that Eq. (2.2) is independent of the heavy fermion mass. For the top quark contribution alone, with $\tan\beta = 1$, one obtains a decay length in the pseudoscalar rest frame of 30 centimeters for $m_A = 100$ MeV. Note that if one considers all quarks and leptons except the first generation fields, then the decay width is increased by a factor of 16, which would correspond to a decay length of 2 centimeters. For $\tan\beta \sim 1$, the branching ratio into photons is 10% for $M_A = 100$ MeV and 40% for $M_A = 200$ MeV. Thus, we see that typical decay lengths, for $\tan\beta = 1$, are on the order of a centimeter. For Model I with small $\cot\beta$, this decay length is increased by a factor of $\tan^2\beta$. These decay lengths will, of course, be increased by a relativistic factor if the pseudoscalar has a large momentum (as it does in $B$-decays).

B. K decays

It is has been long known that the strongest bounds on axion models come from the decay $K \rightarrow \pi A$ [12, 13]; one expects that the same process will significantly constrain the light pseudoscalar scenarios of interest to us here. While many early analyses (that did not
take into account the heaviness of the top quark) seemed to exclude the possibility of a light pseudoscalar in the standard two-doublet scenarios, more recent work suggests that an allowed window remains. It was pointed out by Grzadkowski and Pawelczyk that there are two contributions to the decay amplitude and that the sum may vanish for some choices of model parameters [14]. The first is a direct decay contribution involving the top quark and charged Higgs bosons at one loop; the second is an indirect contribution following from mixing between the axion and the $\pi^0$, $\eta$ and the $\eta'$. We refer the reader to Ref. [14] for the full expressions. As an example, the amplitude for $K^+ \to \pi^+ A$ in Model I can be written schematically as

$$\lambda_w \cot \beta F(m_K, m_{\pi}, m_A, m_\eta, m_{\eta'}) + \cot \beta G(\beta, m_{top}, m_{H^+}, U_{CKM}) .$$

(2.3)

The first term depends only on meson masses and is due to the pseudoscalar mixing; $\lambda_w$ is a chiral Lagrangian parameter that is fixed by the data to be $|\lambda_w| = 3.2 \times 10^{-7}$ [15]. The sign of $\lambda$ can be determined by matching chiral Lagrangian amplitudes to electroweak results [3] and is negative (the imaginary part is proportional to the CP violating factor $\epsilon$ [16] and is thus negligible). The second term represents the direct, one-loop decay amplitude, and depends on the top mass, the charged Higgs mass, and on CKM angles. Specifically, the second term may be written

$$-\frac{1}{2}(m_{\pi}^2 - m_K^2)\frac{\xi}{v}$$

(2.4)

FIG. 1: The branching ratio for $K_L \to \pi^0 A$ for two values of $\tan \beta$ as a function of the charged Higgs mass. We choose $M_A = 150$ MeV. The experimental bound is approximately $4 \times 10^{-8}$.
where

\[ \xi = \frac{G_F}{16\pi^2} \sum_q U_{qs} U_{qd}^* m_q^2 \cot \beta (A_1 + \cot^2 \beta A_2). \]  

(2.5)

Here \( A_1 \) and \( A_2 \) are functions of the top, charged Higgs and W masses and are given explicitly in Ref. [13]. Numerically, the first term of Eq. (2.5) is typically a few times \( 10^{-11} \) GeV and the second is typically \( 10^{-9} \) GeV. However, Grzadkowski and Pawelczyk show that the second term changes sign as the charged Higgs mass varies from 50 GeV to 1000 GeV, and thus at some value the total amplitude vanishes. We have plotted their results for the \( K_L \) decay in Fig. 1, setting \( \tan \beta = 1 \) (so our results then apply to both Model I and Model II), and also \( \tan \beta = 50 \) in Model I. Consideration of \( K^+ \) and \( K_S \) decays leads to qualitatively similar results.

| \( \tan \beta \) | \( K_S \) | \( K_L \) | \( K^\pm \) | \( B \) |
|-----------------|--------|--------|--------|--------|
| 1               | 661-693| 668-672| 669-672| 662-678|
| 2               | 576-643| 597-607| 599-606| 599-605|
| 3               | 546-648| 580-596| 583-594| 584-592|
| 4               | 526-662| 572-594| 576-591| 578-588|
| 5               | 508-679| 567-595| 571-591| 575-587|
| 10              | 434-781| 550-607| 558-599| 566-590|
| 15              | 371-900| 536-621| 548-609| 560-595|
| 20              | 317-1036| 522-637| 538-620| 554-601|
| 30              | 227-1369| 496-669| 518-642| 542-614|
| 40              | 158-1804| 472-702| 500-665| 531-626|
| 50              | 105-2370| 448-738| 482-689| 520-639|

TABLE I: The allowed ranges for the charged Higgs mass (in GeV) for \( K_S, K_L, K^\pm, \) and \( B \) decays. The four ranges overlap for all \( \tan \beta \) shown.

\( \xi \)From Fig. 1 we see that there is a very narrow region of parameter space in which the branching ratio is suppressed. We now must consider whether the experimental bounds on \( K_S, K_L \) and \( K^\pm \) decays can be satisfied simultaneously. The Higgs Hunters Guide [13] refers to two experiments [17, 18] that search for the decay chain \( K^+ \rightarrow \pi^+ A, A \rightarrow e^+ e^- \), and obtain upper limits on the \( \pi A \) branching ratio of order \( 10^{-8} \). However, it is important to
point out that a region between $m_A = 100 - 150$ MeV remains unconstrained due to the large background from the standard decay $K^+ \to \pi^+\pi^0$, followed by $\pi^0$ Dalitz decays. Without precise vertex detection, this can not be distinguished from the pseudoscalar signal. In the particular case of Model I with large $\tan\beta$, the decay length increases by $\tan^2\beta$, and can be several meters. The pseudoscalar would then escape the detector. In that event, bounds from $K^+ \to \pi^+ nothing$ [19, 20, 21], which range from $10^{-7}$ to $10^{-10}$, would apply. Again, the weaker $\mathcal{O}(10^{-7})$ bound applies to a mass interval between $m_A = 130 - 160$ MeV, as a consequence of larger experimental backgrounds. On the other hand, the experimental bounds on the decay $K_L \to \pi^0 A$, are uniformly strong over the entire range of pseudoscalar masses [22]. Fortunately, one can fine-tune the charged Higgs mass to avoid contradiction with both charged or neutral kaon decay bounds. In Table I, we show the required range of charged Higgs masses for $K^+$, $K_L$ and $K_S$ decays. It has been assumed that the $A$ mass is 150 MeV, so that the tighter experimental bounds in charged K decays apply; if the mass is between 100 MeV and 150 MeV, these bounds are relaxed and the ranges for $K^+$ and $K_S$ decays are much wider. For all values of $\tan\beta$ shown in Table I, the allowed ranges for charged Higgs mass overlap and all the bounds can be satisfied with a single fine tuning. For Model II, in which $\tan\beta \sim 1$, the charged Higgs mass must be tuned to approximately one percent precision, but in Model I with larger $\tan\beta$, relatively mild fine-tuning is sufficient. Thus, kaon decays cannot completely exclude the existence of a pseudoscalar in the 100–200 MeV mass range.

C. B decays

In B decays into $KA$, the pseudoscalar will have a relativistic gamma factor of $12 - 24$, depending on its rest mass. Thus, the decay length into electrons will be approximately 25 centimeters (times $\tan^2\beta$). Because of the larger CKM mixing with the top quark, the Higgs-top loop contribution to the amplitude generally dominates over the mixing term by a larger amount than in the case of kaons. A simple estimate illustrates that the branching fraction is potentially large: The loop term involves CKM factors that are comparable to those found in tree-level semileptonic decays, while the $16\pi^2$ in the loop is partly compensated by the smaller two-body phase space. The resulting prediction has a shape very similar to that
for $K$ decays in Fig. 1. Again, there is a narrow region of parameter-space where the rate vanishes, and this region matches the narrow region in $K$-decays. This is not surprising since the analog of Eq. (2.5) for $B$ decays has the same functional dependence on the charged Higgs mass, up to an overall factor. One might hope that higher order corrections would separate the $K$ and $B$ decay allowed mass windows, but a one-percent effect would not be sufficient to alter our qualitative results.

What are the experimental limits? Recently, the BELLE Collaboration published a value for the branching fraction for $B \rightarrow K e^+ e^-$ of $0.75 \pm 0.2 \times 10^{-6}$ [23]. Since this is in agreement with theory, a bound on new physics contributions of approximately $2 \times 10^{-7}$ can be obtained. However, the BELLE analysis included a mass cut on the electron-positron pair of 140 MeV, to suppress background from photon conversions and $\pi^0$ Dalitz decays. Thus, the bound does not apply to the 100–140 MeV window. The CLEO Collaboration has searched for $B^\pm \rightarrow K^\pm$ nothing and $B^0 \rightarrow K^0_S$ nothing decays, and obtains a bound on the branching ratios of $5 \times 10^{-5}$ [24]. While this does cover the mass range in which the BELLE analysis does not apply, it is only relevant if all the pseudoscalars escape detection. For masses between 100 and 140 MeV, one can ask what fraction of the $A$’s will escape the detector. For $\tan \beta = 50$, the decay length will be over 10 meters and almost all of the $A$’s would escape; the CLEO bound would then apply. In general, approximately $e^{-4 \cot^2 \beta}$ of the $A$’s escape the detector, which is a barrel calorimeter of roughly a meter radius. The bound would then be weaker by this factor, or $5 \times 10^{-5} e^{4 \cot^2 \beta}$ for the branching ratio. Using this experimental bound, we find the allowed charged Higgs mass range given in Table I. We see that the same fine-tuning needed (for $\tan \beta \sim 1$) for kaon decays will automatically suppress the B-decay rate.

We conclude that neither model I nor II can be definitively excluded from the bounds from $B$ decay, although fine-tuning is needed if $\tan \beta \sim 1$, as required in Model II.

**D. Leptophobic Pseudoscalars**

As noted in the introduction, it is simple to have a three Higgs model in which two of the Higgs doublets couple to quarks (with Model I or Model II couplings) and a third couples to leptons. If the third doublet does not mix with the others, the leptonic couplings of the light
pseudoscalar are eliminated. The $K$ and $B$ decays discussed in the previous two subsections will generally not be affected in such a model. However, the decay of the pseudoscalar will now be entirely into photon pairs and the lifetime will generally be 2-3 times larger than the usual case. Note that in the 100 − 140 GeV mass window, the stronger bounds from $K_L$ decays and from CLEO will certainly apply (without any significant exponential correction for decays inside the detector). Again, these bounds can be evaded with a suitable fine tuning of the charged Higgs mass.

III. FERMIOPHOBIC PSEUDOSCALARS

We have seen in the previous sections that a pseudoscalar state in the 100 to 200 MeV mass range is consistent with the stringent bounds from $K$ and $B$ meson decays. However, in the conventional scenarios considered thus far, this result follows from an accidental zero in the decay amplitudes, as well as a willingness to accept fine tuning. In this section we consider another possibility, that the couplings of the pseudoscalar to matter are naturally suppressed. After discussing the experimental bounds, we argue that a natural place to search for such a state is in a low-energy photoproduction experiment, such as those possible at Jefferson Lab. We estimate the production rate and comment on the relevant discovery signal in Section IV.

We have already stated the motivation for considering a pseudoscalar state that is light: it might be the would-be goldstone boson associated with a global symmetry that is only approximate. In the introduction, we outlined a plausible scenario with a singlet scalar and a vectorlike multiplet in a complete $SU(5)$ representation, taken to be a $5 + \bar{5}$ for simplicity.

Since the exotic matter is vector-like, it can be made arbitrarily heavy and integrated out of the theory. This leads to nonrenormalizable interactions between the pseudoscalar and the standard model gauge fields. If $M_F$ is the mass scale of the vector-like matter $\psi$, and the pseudoscalar coupling is given by $(iA\lambda/\sqrt{2})\bar{\psi}\gamma^5\psi$, then one obtains

$$\mathcal{L} = \frac{q^2\lambda}{32\sqrt{2}\pi^2 M_F} \epsilon^{\mu\nu\rho\sigma} AF_{\mu\nu} F_{\rho\sigma}$$

for the effective coupling of the pseudoscalar to two photons. Here $q$ represents the electric charge of $\psi$, and $F_{\mu\nu}$ is the electromagnetic field strength. Note that this can be generalized
to any non-Abelian gauge group by replacing $q^2$ with the Casimir $T_F$ (defined by $\text{Tr}[T^a T^b] = T_F \delta^{ab}$) and by summing over the field strength tensors. For a fermiophobic pseudoscalar in the mass range of interest to us, the only possible decay is to two photons, and from Eq. (3.1) we obtain the decay width

$$\Gamma(A \to \gamma\gamma) = \frac{16}{9} \cdot \frac{\alpha^2 \lambda^2 m_A^3}{128 \pi^3 M_F^2}.$$ (3.2)

If $M_F$ is not far above the top quark mass, say 200 GeV, and $\lambda = 1$, then one obtains a lifetime

$$\tau(M_F = 200 \text{ GeV}) = 1.1 \times 10^{-3} \text{ sec} \left(\frac{\text{MeV}}{m_a}\right)^3.$$ (3.3)

For energies of a few GeV, typical of the photoproduction experiments that we will mention later, the pseudoscalar can travel a macroscopic distance before it decays. A pseudoscalar with a mass of 150 MeV and an energy of 3 GeV will have a decay distance of 160 centimeters.

One might think that the scenario described above is relatively insensitive to the bounds from meson decays due to the weakness of the pseudoscalar’s coupling to ordinary matter. However, the experimental bounds on the branching fraction of $K$ or $B$ mesons to $\pi$ + pseudoscalar are so stringent that operators like Eq. (3.1) are potentially significant, even when they contribute only at one loop. Here we estimate the contribution to $K \to \pi A$ in order to constrain the parameter space of the model. We comment on the constraints from $B$ decays at the end of this section.

The operator with the largest potential effect on low-energy hadronic decays is the gluonic version of Eq. (3.1). We use a chiral lagrangian approach to estimate the branching fraction of interest [14]. First we represent the light pseudoscalar nonet via the nonlinear representation

$$\Sigma = \exp(2i\pi/f_\pi)$$ (3.4)

where $f_\pi = 93$ MeV is the pion decay constant, and where $\pi$ is the matrix of fields

$$\pi = \left( \begin{array}{ccc}
\frac{\pi^0}{\sqrt{2}} + \frac{\eta}{2\sqrt{3}} + \frac{\eta'}{\sqrt{6}} & \frac{\pi^+}{\sqrt{2}} & K^+ \\
-\frac{\pi^-}{\sqrt{2}} + \frac{\eta}{2\sqrt{3}} + \frac{\eta'}{\sqrt{6}} & \frac{\pi^0}{\sqrt{2}} + \frac{\pi^+}{\sqrt{2}} + \frac{\eta}{\sqrt{3}} & \frac{1}{2}(K^0_s + K^0_L) \\
\frac{K^-}{\sqrt{2}} & -\frac{1}{2}(K^0_s - K^0_L) & -\frac{\eta}{\sqrt{6}} + \frac{\eta'}{\sqrt{6}} \end{array} \right)$$ (3.5)

Here we have ignored CP violation and expressed the neutral kaons in terms of their CP eigenstates. Also note that we have chosen to include the $\eta'$, so that $\Sigma$ is an element of $U(3)$
rather than SU(3). The Σ field transforms simply under the chiral SU(3) symmetry

$$\Sigma \rightarrow U_L^\dagger \Sigma U_R$$

leading to the usual lowest order effective Lagrangian

$$\mathcal{L}_0 = \frac{f^2}{4} \text{Tr} \partial_\mu \Sigma^\dagger \partial^\mu \Sigma + \frac{1}{2} f^2 \mu \text{Tr} (M\Sigma^\dagger + \Sigma M^\dagger) ,$$

where $M$ represents the light quark current mass matrix. However, Eq. (3.7) does not take into account the QCD anomaly, which relates the divergence of the axial current to the product of gluon field strength tensors $G^{\mu\nu} \tilde{G}_{\mu\nu}$. A possible method of incorporating this effect into the chiral lagrangian is to introduce the additional terms [25]

$$\mathcal{L}_{\text{anom}} = \frac{1}{2} i q(x) \log \frac{\det \Sigma}{\det \Sigma^\dagger} + cq(x)^2$$

where $q(x)$ represents

$$q(x) = \frac{g^2}{32\pi^2} G^a_{\mu\nu} \tilde{G}^{\mu\nu}_a .$$

Under an axial U(1) rotation, the field $\Sigma$ is multiplied by an overall phase, and it is not hard to show that $\mathcal{L}_0 + \mathcal{L}_{\text{anom}}$ yields the appropriate divergence of the axial vector current [25]. Now, one may treat $q(x)$ as an auxiliary “glueball” field, and remove it using its equation of motion. One then finds

$$\mathcal{L}_{\text{anom}} = -\frac{1}{4c} \left( \frac{i}{2} \log \frac{\det \Sigma}{\det \Sigma^\dagger} \right)^2 .$$

This term determines the $\eta'$ mass, and the parameter $c$ can be chosen accordingly. If one now includes the pseudoscalar coupling to gluons, an additional term must be added to Eq. (3.8), namely $A q(x)/(2\sqrt{2}M_F)$, in which case Eq. (3.10) is modified

$$\mathcal{L}_{\text{anom}} = -\frac{1}{4c} \left( \frac{i}{2} \log \frac{\det \Sigma}{\det \Sigma^\dagger} + \frac{1}{2\sqrt{2}M_F} A \right)^2 .$$

This interaction leads to mass mixing between the pseudoscalar and the $\eta'$; we find that the mixing angle is given approximately by

$$\theta_{\Delta S} \approx \frac{1}{4\sqrt{3}} \frac{f_\pi}{M_F}$$

or numerically, $7 \times 10^{-5} \cdot (200 \text{ GeV}/M_F)$. We may extract the $\Delta S = 1$ $K\pi\eta'$ vertex from the chiral Lagrangian term

$$\mathcal{L}_{\Delta S=1} = \frac{f^2}{4} \text{Tr} (\lambda_w h \partial_\mu \Sigma \partial^\mu \Sigma^\dagger)$$

(3.13)
where \( h \) is octet-dominant \( \Delta S = 1 \) spurion

\[
h = \begin{pmatrix}
0 & 0 & 0 \\
0 & 0 & 1 \\
0 & 0 & 0
\end{pmatrix},
\tag{3.14}
\]

and \( \lambda_w = 3.2 \times 10^{-7} \) is a parameter that takes into account the strength of the weak interactions \[15\]. We find

\[
\Gamma(K^+ \to \pi^+ A) = \frac{1}{384\pi} \frac{\lambda_w^2 \theta_{A'0}^2}{m_A^3 f_\pi^2} \left( m_A^2 + 2m_K^2 \right)^2 \left[ (m_K^2 - m_\pi^2 + m_A^2)^2 - 4m_K^2 m_A^2 \right]^{1/2}. \tag{3.15}
\]

As a point of reference, if one sets \( m_A = 100 \text{ MeV} \), one obtains the branching fraction \( 5.6 \times 10^{-7} \cdot (200 \text{ GeV}/M)^2 \).

Different experimental bounds are relevant depending on the lifetime and boost of the pseudoscalar. If the pseudoscalar decays inside the experimental detector, the relevant bound on the \( K^+ \) branching fraction is \[3\]

\[
BF(K^+ \to \pi^+ \gamma \gamma) < 1.4 \times 10^{-6}. \tag{3.16}
\]

If the pseudoscalar escapes the detector unobserved, one must contend with more stringent bounds, ranging from \( \sim 10^{-7} \) to \( \sim 10^{-10} \), depending on the pseudoscalar mass \[21\]. In Fig. 2

![FIG. 2: Allowed parameter space for the fermiophobic scenario.](image)

we display the allowed region of the model’s parameter space. Within the two excluded
regions toward the top of the figure, the pseudoscalar is long lived enough to escape the
detector, while the branching fraction exceeds the bounds given in Ref. [21]. The gap
between these regions corresponds to a mass interval in which there are larger experimental
backgrounds. Immediately below each of these excluded regions, the pseudoscalar decays to
two photons within the detector (assumed to have a fiducial length scale of 1.45 meters [26])
and the weaker bound in Eq. (3.16) becomes relevant. However, one never reaches the region
of parameter space excluded by the $K^+ \rightarrow \pi^+ \gamma \gamma$ bound since the vector-like matter would
itself become light enough to be detected in direct collider searches. We will restrict ourselves
to the allowed regions of Fig. 2 with smallest $M_F$ in discussing pseudoscalar production rates,
in the next section.

Finally, we should comment on the bounds from the analogous decays of neutral kaons
and $B$ mesons. First, the $K^0_s$ indeed may decay into $\pi^0 A$; however, the total width of the
$K^0_s$ is approximately two orders of magnitude larger that that of the $K^+$, so the branching
fraction to the decay mode of interest is suppressed by this factor relative to our previous
results. We therefore obtain no further bounds. The $K^0_L$, on the other hand, has a total
width that is about a factor of four smaller than that of the charged kaon. However, the
decay $K^0_L \rightarrow \pi^0 A$ is CP violating, so that the decay amplitude is suppressed by an additional
CP-violating spurion factor of $\sim 10^{-3}$ [16], and again no further bound is obtained. In the
$B$ system, the decay $B \rightarrow K \eta'$ is observed, and has a branching fraction of order $10^{-5}$ [27].
Using our previous result for the $A \eta'$ mixing angle, we estimate that the branching fraction
for $B \rightarrow KA$ is $O(10^{-15})$ and no further bound is obtained.

IV. PRODUCTION AT JEFFERSON LAB

We have seen that there is a window for light pseudoscalars in the $100 - 200$ MeV mass
range. For the two-doublet Model II (or Model I with $\tan \beta \sim 1$) the window requires
substantial fine-tuning of the charged Higgs mass; for the two-doublet Model I with large
$\tan \beta$, there is less fine-tuning, and for the fermiophobic case there is a very large region of
allowed parameter space. How can one detect these pseudoscalars?

A number of authors have considered light pseudoscalar detection at high-energy colliders [10, 28]. Larios, Tavares-Velasco and Yuan [10] discussed production at the Tevatron,
the LHC and future colliders. They focused on the two-photon decay mode, which at high energies registers as a single photon signature. In this section, we consider the possibility of detecting the pseudoscalars we have discussed in a beam dump experiment at the Thomas Jefferson National Accelerator Facility (Jefferson Lab).

Jefferson Lab has a high intensity photon beam directed into the CLAS detector in Hall B. The maximum energy is currently 6 GeV with an upgrade to 12 GeV planned. The photon beam has a bremsstrahlung spectrum with a luminosity of approximately $10^{34}$ cm$^{-2}$ sec$^{-1}$ if the photons are untagged. At the 12 GeV upgrade a monochromatic 9 GeV photon beam will also be available, with a luminosity of approximately $10^{33}$ cm$^{-2}$ sec$^{-1}$. The amplitude for pseudoscalar photoproduction may receive two possible contributions. In the conventional two-Higgs-doublet models, one can photoproduce the pseudoscalar most copiously off the strange quark sea in the proton. Second, in all the models we have discussed, the pseudoscalar may bremsstrahlung off the incident photon via the loop-induced $A\gamma\gamma$ vertex.

Once produced, the pseudoscalar will travel some distance and then decay into either $e^+e^-$ or $\gamma\gamma$, depending on the model. If the beam dump consists of a meter or more of material, then most of the $\gamma\gamma$ background events will be suppressed. It is thus important that the lifetime of the $A$ be sufficiently long that a substantial number make it through the beam dump.

We first concentrate on production. Consider photoproduction of the pseudoscalar off the strange quark in the proton. The parton level cross section in the center of mass frame is

$$\frac{d\hat{\sigma}}{d\cos\theta} = \frac{h^2 e^2 p}{144\pi \hat{s}^{3/2}} \left[ \frac{m_s^2 - \hat{t}}{\hat{s} - m_s^2} + \frac{2m_s^2 m_A^2}{(\hat{s} - m_s^2)^2} + \frac{\hat{s} - m_s^2}{m_s^2 - \hat{t}} \right. $$

$$ + \frac{2m_s^2 m_A^2}{(m_s^2 - \hat{t})^2} + \frac{\hat{t}\hat{s} - (m_A^2 + m_s^2)(\hat{s} + \hat{t}) + m_s^4 + m_A^4}{(\hat{s} - m_s^2)(m_s^2 - \hat{t})} \left. \right] .$$

(4.1)

Here, $h$ is the Yukawa coupling of the $A$ to the strange quark, $p$ is the $A$ momentum; we have approximated the initial photon momentum as $\sqrt{\hat{s}}/2$ in the phase space factors to simplify the expression. In finding the full cross section for photoproduction, we multiply by the parton distribution function for the strange quark and integrate. However, since the parton model becomes less reliable at small momentum transfers, one must keep in mind that there is significant theoretical uncertainty from the small $x$ region of integration, where the partonic cross section is largest. We therefore cut off the $x$ integration at a value where
\( \hat{s} = xs = 1 \text{ GeV}^2 \). We believe that this choice is reasonable. At lower \( \hat{s} \) there will not be enough energy to produce \( \phi, \eta \) and \( K \) mesons, and thus one expects an additional suppression from the electromagnetic form factor due to the decrease in available exclusive channels. The resulting cross section is rather insensitive to the beam energy, varying from 3.6 cot\(^2\) \( \beta \) to 2.0 cot\(^2\) \( \beta \) femtobarns as the photon energy varies from 2 to 12 GeV. For a luminosity of \( 10^{34} \text{ cm}^{-2} \text{ sec}^{-1} \), this will yield approximately 800 cot\(^2\) \( \beta \) events per year. In order to be detected, these pseudoscalars must travel through a beam dump. The lifetime, as discussed in Section II, gives a decay length for a 100 MeV pseudoscalar of 1.2 tan\(^2\) \( \beta \) centimeters times the relativistic factor of \( E/M_A \). Consider a 6 GeV beam and tan \( \beta = 1 \). The decay length is then 72 centimeters, and roughly 25\% of the particles, or 200 particles/year, will travel through a one-meter beam dump. Since the differential cross section has a \( t \)-channel pole in the massless quark limit, it is forward peaked and this estimate will not suffer a substantial solid angle dilution. As tan \( \beta \) increases, the production cross section drops, but the decay length increases. In Table II, we show the number of events that traverse a one-meter beam dump per year, assuming \( 10^{34} \text{ cm}^{-2} \text{ sec}^{-1} \) luminosity. These pseudoscalars will primarily decay into an electron-positron pair. One should keep in mind that the uncertainties caused by the low \( x \) cutoff could be substantial, and thus these event rates are approximate. Also, for larger beam energies and larger tan \( \beta \), the decay length will be too long for a substantial number of events to occur in a detector. Nonetheless, the relatively high event rate indicates that further experimental analysis is warranted.

The second production mechanism is through the \( A\gamma\gamma \) vertex. In the two-Higgs-doublet models, the production mechanism already considered strongly dominates, but in the fermiophobic model, pseudoscalar bremsstrahlung off the incident photon is the only possibility. The parton level cross section is

\[
\frac{d\hat{\sigma}}{d\cos \theta} = -\frac{\lambda^2 Q^2 e^6 \left(p(2m_q^2m_A^4 + \hat{t}^3 - 2(m_A^2 - \hat{s})\hat{t}^2 + ((m_A^2 - m_q^2)^2 + m_q^4 + 2\hat{s}^2 - 2(m_A^2 + 2m_q^2)\hat{s})\hat{t})\right)}{2048\pi^5 M_F^2 \hat{t}^2 \hat{s}^{3/2}} \tag{4.2}
\]

where \( \lambda \) is the coupling of the fermion in the loop to the \( A \), \( M_F \) is the mass of the fermion in the loop, \( p \) is the final state 3-momentum of the \( A \), and \( Q \) is the quark charge in units of e. In deriving this expression, we have assumed that \( M_F \) is much greater than the photon energy (certainly true for the fermiophobic case). For \( \sqrt{\hat{s}} \gg m_q + m_A \), we find that \( \hat{\sigma} \) is
\[
\tan \beta \quad 4 \text{ GeV} \quad 6 \text{ GeV} \quad 9 \text{ GeV} \quad 12 \text{ GeV} \quad 24 \text{ GeV}
\]
\[
\begin{array}{cccccc}
1 & 125 & 210 & 350 & 340 & 340 \\
2 & 150 & 147 & 144 & 135 & 120 \\
3 & 88 & 85 & 77 & 66 & 46 \\
4 & 53 & 53 & 48 & 40 & 32 \\
5 & 41 & 34 & 28 & 25 & 21 \\
\end{array}
\]

TABLE II: The number of pseudoscalars traversing at least one meter for various values of the beam energy and \( \tan \beta \) in the two-doublet model with photoproduction off the strange quark sea. We have assumed a luminosity of \( 10^{34} \text{ cm}^{-2} \text{ sec}^{-1} \) and a pseudoscalar mass of 100 MeV. Most will decay into an electron-positron pair.

well approximated by

\[
\hat{\sigma} \approx \frac{\alpha^3 Q^2 \lambda^2}{64 \pi^2 M_F^2} \frac{1}{s} \left[ (2\hat{s} - m_A^2)^2 \log \frac{(\hat{s} - m_A^2)^2}{m_q m_A^2} - 3(\hat{s} - m_A^2)^2 \right].
\] (4.3)

The exact parton-level total cross section is shown in Fig. 3. The approximate expression given in Eq. (4.3) yields results that are visually indistinguishable from those shown in Fig. 3. Using CTEQ set 5L structure functions for the up and down sea and valence quarks

FIG. 3: Parton-level production photoproduction cross section in the fermiophobic scenario as a function of center of mass energy, with \( M_F = 200 \text{ GeV} \) and \( \lambda = 1. \)
we obtain the total production cross section shown in Fig. 4. Assuming a monochromatic photon beam and a Jlab-like luminosity of $10^{34} \text{ cm}^{-2} \text{s}^{-1}$, one estimates 315 production events per year per femtobarn of total cross section; qualitatively speaking, Fig. 4 suggests $\mathcal{O}(10^2)$ events per year at an energy-upgraded Jlab, or at some similar facility.

![Graph showing photoproduction cross section](image)

FIG. 4: Photoproduction cross section in the fermiophobic scenario, as a function of photon beam energy in the lab frame, with $M_F = 200$ GeV and $\lambda = 1$.

A more realistic analysis would take into account that the highest luminosity photon beam at Jlab is not monoenergetic, but has a bremsstrahlung spectrum. We approximate this effect by assuming a total luminosity of $10^{34} \text{ cm}^{-2} \text{s}^{-1}$, with a distribution $dL/dE_\gamma \propto 1/E_\gamma$, with $E_\gamma$ ranging from 1 GeV up to the beam energy; the event rate is determined by the integral

$$\int \sigma \frac{dL}{dE_\gamma} dE_\gamma . \quad (4.4)$$

Table III shows the events per year for a number of different choices for the beam energy and pseudoscalar mass. Unlike the two-doublet model, the lifetime discussed in Section III is sufficiently long that most of these pseudoscalars will traverse a one-meter beam dump. Another major difference is that these pseudoscalars will decay into two photons, i.e. they will look like long-lived $\pi^0$'s. A more detailed analysis taking into account possible experimental acceptances and cuts would be needed to determine whether this signal could be separated from background under realistic conditions.
| $E_\gamma$ (GeV) | $m_A = 100$ MeV | $m_A = 200$ MeV |
|-----------------|-----------------|-----------------|
| 6               | 99              | 79              |
| 12              | 108             | 87              |
| 24              | 117             | 96              |

TABLE III: Photoproduction event rate per year in the fermiophobic scenario, with $M_F = 200$ GeV and $\lambda = 1$. The total luminosity is taken to be $10^{34}$ cm$^{-2}$ s$^{-1}$ and a Bremstrahlung photon spectrum is assumed between 1 GeV and the beam energy.

V. CONCLUSIONS

We have considered light, elementary pseudoscalars with masses between 100 and 200 MeV. We have argued that such states may evade the stringent bounds from $K$ and $B$ meson decays, while remaining of interest in searches at low-energy photoproduction experiments, such as those possible at Jefferson Lab. In conventional two-Higgs doublet models, light pseudoscalars may evade the strange and bottom meson decay bounds due to a possible cancellation in the decay amplitude. In this case, the coupling of the pseudoscalar to quarks is substantial and one can produce the pseudoscalar state copiously via photoproduction off the strange quark sea in a nucleon target. On other hand, if one wishes to avoid fine tuning in evading the decay bounds, one can consider very natural scenarios in which the pseudoscalar is fermiophobic. We have presented one concrete realization of this idea, motivated by extra dimensions, and have isolated the allowed parameter space of the model. In the fermiophobic scenario, the pseudoscalar-two photon coupling leads to production via pseudoscalar bremsstrahlung off the incoming photon line. The event rate is substantial enough to make accelerator searches of potential interest.

Acknowledgments

We thank Andrew Bazarko, Morton Eckhause, Keith Griffioen, Bohdan Grzadkowski, Jon Urheim and C.P. Yuan for useful comments. We thanks the National Science Foundation (NSF) for support under Grant No. PHY-9900657. In addition, C.D.C. thanks the NSF for
support under Grant Nos. PHY-0140012 and PHY-0243768.

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