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Constraints on a very light CP–odd Higgs of the NMSSM and other axion–like particles

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Abstract

In the NMSSM, a light CP–odd Higgs arises due to spontaneous breaking of approximate symmetries such as Peccei–Quinn or R–symmetry and is motivated by string theory. The case when it is heavier than two muons is well studied and constrained. We analyze various meson decay, $g-2$, beam dump and reactor bounds on the CP–odd Higgs with mass below the muon threshold, in particular, addressing the question how light a CP–odd Higgs can be. We find that it has to be heavier than 210 MeV or have couplings to fermions 4 orders of magnitude below those of the Standard Model Higgs. Our analysis applies more generally to couplings of a light pseudoscalar to matter.

Contents

1 Introduction **2**

2 NMSSM and a light CP–odd Higgs **3**
2.1 Peccei–Quinn limit 3
2.2 R–symmetry limit 4

3 Constraints on a light CP–odd Higgs **4**
3.1 Rare B–decays $B \to K +$ invisible 7
3.2 Rare K–decays $K \to \pi +$ invisible 7
3.3 Rare decays $B \to K\ e^+e^-$, $K \to \pi\ e^+e^-$ 8
3.4 Rare K–decays $K^+ \to \pi^+ + X$ 10
3.5 Radiative Upsilon–decays 10
3.6 Pion decay $\pi^0 \to e^+e^-$ 11
3.7 Muon anomalous magnetic moment 12
3.8 Other constraints 12

4 Further bounds from reactor and beam dump experiments **13**
4.1 Reactor bounds 13
4.2 Beam dump limits 14

5 Conclusion **16**

References **21**

1 Introduction

An attractive extension of the minimal supersymmetric Standard Model (MSSM) is the next–to–minimal supersymmetric Standard Model (NMSSM). Historically, it has been motivated by the μ–problem of the MSSM which can be solved by replacing the μ–parameter with an SM singlet S [1]. More recently, it has been suggested that the fine–tuning problem of the MSSM can be alleviated or removed in the NMSSM when a light CP–odd Higgs is present in the spectrum [2, 3]. Although this scenario is now tightly constrained by the new ALEPH analysis [4] as well as BaBar data on $\Upsilon(3S)$ decays [5, 6], some parameter space remains available [7] and can further be probed by η_b–decays [8].

The NMSSM–like models can also be obtained from the heterotic string [9], based on the MSSM constructions of Refs. [10, 11, 12, 13]. An interesting feature of these models is that the SM singlet S is not a singlet under the full gauge symmetry of the heterotic string ($E_8 \times E_8$). Therefore all self–interactions

$$S\ ,\ S^2\ ,\ S^3\ ,\ ... \tag{1}$$

are forbidden by gauge invariance. They are only allowed after spontaneous breaking of gauge symmetry and the corresponding couplings are suppressed by VEVs of the symmetry–breaking fields. On the other hand, interactions like SH_1H_2 can be allowed without symmetry breaking.

Due to this hierarchy in the couplings, one obtains specific versions of the NMSSM such as the Peccei–Quinn (PQ) or decoupling versions [9]. In the former case, a pseudo–Goldstone boson appears in the spectrum at the electroweak scale. Its mass is generated by small PQ violating effects and can be much below the GeV scale.

Motivated by these considerations and also by a possible connection of a light pseudoscalar to dark matter [14, 15, 16, 17], in this work we study constraints on a very light CP–odd Higgs and address the question how light a CP–odd Higgs can be. We analyze bounds from various meson decays, muon $g-2$, beam dump and reactor experiments for the pseudoscalar mass below the muon threshold. Our results apply beyond the NMSSM to couplings of a light pseudoscalar to matter.

2 NMSSM and a light CP–odd Higgs

The NMSSM is the MSSM extended by a singlet superfield S. In what follows, we will focus on a particular version of the NMSSM which has no direct μ–term, the so called Z_3–symmetric NMSSM. The relevant superpotential of the Z_3–symmetric NMSSM is

$$W = \lambda S H_1 H_2 + \frac{1}{3}\kappa S^3 \,, \tag{2}$$

while the soft terms are given by

$$V_{\rm soft} = m_1^2|H_1|^2 + m_2^2|H_2|^2 + m_S^2|S|^2 + \left(\lambda A_\lambda S H_1 H_2 + \frac{1}{3}\kappa A_\kappa S^3 + \text{h.c.}\right) . \tag{3}$$

A light pseudoscalar A^0 appears naturally in two limiting cases: when the Higgs potential possesses either approximate Peccei–Quinn (PQ) or approximate R–symmetry [18, 19, 20, 21, 3, 22, 23].

2.1 Peccei–Quinn limit

In the limit $\kappa \to 0$, the Lagrangian is invariant under the transformation

$$H_{1,2} \to e^{i\alpha} H_{1,2} \quad , \quad S \to e^{-2i\alpha} S \,. \tag{4}$$

At the electroweak scale this symmetry gets broken, resulting in the "axion" [24]

$$\begin{aligned} A^0 &= \frac{1}{N}\Big(v\ \sin 2\beta\ A^0_{\rm MSSM} - 2s\ S_I\Big) \,, \\ N &= \sqrt{v^2 \sin^2 2\beta + 4s^2} \,, \end{aligned} \tag{5}$$

where $A^0_{\rm MSSM} = \cos\beta\ H_{1I} + \sin\beta\ H_{2I}$ is the MSSM pseudoscalar, $s = \langle S \rangle$ and the subscript I refers to the imaginary part of the Higgs neutral component. As usual, $\tan\beta = v_1/v_2$ and $v = \sqrt{v_1^2 + v_2^2} = 174$ GeV. The mass of the pseudoscalar is most easily expressed in the large $\tan\beta$ regime [19]:

$$m^2_{A^0} \simeq -3\kappa A_\kappa s \,. \tag{6}$$

Since renormalization of κ is proportional to κ itself, this coupling can be very small and A^0 very light. In the string NMSSM example of Ref. [9], $\kappa < \mathcal{O}((\phi/M_{\rm Pl})^5)$ with ϕ being an average VEV of certain SM singlets. For ϕ an order of magnitude below the Planck scale, κ can be as small as 10^{-6} leading to a 100 MeV pseudoscalar. In other models it can be even lighter. Since the PQ symmetry is anomalous as in the DFSZ construction [25], the lower limit on m_{A^0} is set by the anomaly contribution of order 100 keV [25] (for $s \sim v$).

2.2 R–symmetry limit

In the limit $A_\kappa, A_\lambda \to 0$, the Higgs sector of the NMSSM is R–invariant. Under R–symmetry the superfields transform as

$$H_{1,2} \to e^{i\alpha_R} H_{1,2} \quad , \quad S \to e^{i\alpha_R} S \; , \tag{7}$$

such that the superpotential transforms with charge 2. Spontaneous breaking of this symmetry results in an "R–axion". Its composition is given by [24]

$$\begin{aligned} A^0 &= \frac{1}{N}\Big(v \;\sin 2\beta \; A^0_{\rm MSSM} + s \; S_I \Big) \; , \\ N &= \sqrt{v^2 \sin^2 2\beta + s^2} \; , \end{aligned} \tag{8}$$

with $A^0_{\rm MSSM}$ as in Eq. 5.

Unlike the Peccei-Quinn symmetry, R–symmetry is not a (classical) symmetry of the full Lagrangian. Even if $A_\kappa, A_\lambda \to 0$, the gaugino mass terms break it explicitly. Non–zero A–terms are induced by renormalization, so their minimal value is a loop factor times the gaugino mass. The axion mass is again approximated by (6).

In both PQ– and R–symmetric cases, the light pseudoscalar is mostly a singlet in the limit $s \gg v \;\sin 2\beta$. Its couplings to gauge bosons and SM matter are suppressed in this limit. However, s cannot be too large, otherwise a large effective μ–term is induced. An exception is the case $\lambda \ll 1$, which corresponds to the "decoupling limit", i.e. no communication between the singlet and the rest of the NMSSM.

3 Constraints on a light CP–odd Higgs

Following the notation of Ref. [7], the coupling of the CP-odd Higgs A^0 to fermions is given by

$$\Delta\mathcal{L} \;=\; -i\frac{g}{2m_W}\; C_{Aff} \left(m_d \; \bar{d}\gamma_5 d + \frac{1}{\tan^2\beta} m_u \; \bar{u}\gamma_5 u + m_l \; \bar{l}\gamma_5 l \right) A^0. \tag{9}$$

In the NMSSM, the coupling C_{Aff} can be expressed in terms of the singlet–doublet mixing angle θ_A and $\tan\beta$: $C_{Aff} = \cos\theta_A \tan\beta$ [7], with $\cos\theta_A = v \;\sin 2\beta/N$ and N given in Eqs. 5,8. In what follows, we treat it as a free parameter and derive various particle physics constraints on it. In the NMSSM, very large ($> 10^2$) and very small ($< 10^{-2}$) values of C_{Aff} lead to violation of perturbativity and/or finetuning, so it usually suffices to focus on the

moderate C_{Aff} window. However, our analysis applies more generally to the coupling of any pseudoscalar to matter as long as the C_{Aff} parameter is universal for all fermions.

For $m_{A^0} > 2m_\mu$, the resulting constraints, in particular from meson decays, are well studied. In the range $2m_\mu < m_{A^0} < 3m_\pi$, A^0 decays predominantly into 2 muons, which is constrained by $K^+ \to \pi^+ A^0$ and $B \to K\ A^0$. The resulting bound is of order $C_{Aff} < \mathcal{O}(10^{-2})$ [26]. Above the 3–pion threshold, the branching ratio for the decays into muons reduces and the bound weakens somewhat. For even larger A^0 mass, the $\Upsilon \to \gamma A^0$ decay imposes $C_{Aff} < 0.5$ [7] at $\tan\beta \sim 1$ with the bound getting weaker, $\mathcal{O}(1)$, close to m_Υ. Above 12 GeV, the DELPHI data on $e^+e^- \to b\bar{b}A^0 \to b\bar{b}b\bar{b}$ set a limit $C_{Aff} < \mathcal{O}(10)$ [7]. Further constraints, usually relevant at large $\tan\beta$, are summarized in Ref. [24].

The $m_{A^0} < 2m_\mu$ territory is less well explored. Some constraints have been studied in Refs. [26, 20] in the framework of the NMSSM, and in Refs. [27, 28] for 2 Higgs Doublet Models (see also [29]). In what follows, we delineate the $\{m_{A^0}, C_{Aff}\}$ parameter space taking into account meson decay, muon $g-2$, nuclear reactor and beam dump constraints. In particular, we address the question how light a CP–odd Higgs boson can be. Since we work in terms of the coupling C_{Aff}, most of our results are largely independent of $\tan\beta$.

For m_{A^0} below the muon threshold, A^0 can only decay into electron-positron pairs and photons. Its total decay width is

$$\Gamma_{\rm tot} = \Gamma(A^0 \to e^+e^-) + \Gamma(A^0 \to \gamma\gamma) \tag{10}$$

with

$$\Gamma(A^0 \to f\bar{f}) = \frac{\sqrt{2}G_F}{8\pi}\ m_f^2\ m_{A^0}\ C_{Aff}^2\ \sqrt{1-4\frac{m_f^2}{m_{A^0}^2}}\ , \tag{11}$$

$$\Gamma(A^0 \to \gamma\gamma) = \frac{\sqrt{2}G_F\alpha^2}{16\pi^3}\ m_{A^0}^3\ \left|\sum_i r C_{Aii} Q_i^2 k_i F(k_i)\right|^2\ , \tag{12}$$

where $r = 1(N_c)$ for leptons (quarks), $k_i = m_i^2/m_{A^0}^2$, Q_i is the charge of the fermion in the loop; $C_{Aii} = C_{Aff}$ for the down–type fermions and $C_{Aii} = C_{Aff}/\tan^2\beta$ for the up–type fermions. (Here we neglect the chargino contribution as the coupling is dominated by the SM fermions.) The loop function $F(k_i)$ is given by [30]

$$F(k_i) = \begin{cases} -2(\arcsin\frac{1}{2\sqrt{k_i}})^2 & \text{for } k_i \geq \frac{1}{4}, \\ \frac{1}{2}[\ln(\frac{1+\sqrt{1-4k_i}}{1-\sqrt{1-4k_i}}) + i\pi]^2 & \text{for } k_i < \frac{1}{4}, \end{cases} \tag{13}$$

and has the limits

$$k_i F(k_i) = \begin{cases} 0 & \text{for } k_i \ll 1, \\ -\frac{\pi^2}{8} & \text{for } k_i = \frac{1}{4}, \\ -\frac{1}{2} & \text{for } k_i \gg 1. \end{cases} \tag{14}$$

For example, with $m_{A^0} = 0.5$ MeV and $C_{Aff} = 1$ the total width is about 4×10^{-12} eV. This corresponds to the decay length (for a boost factor $\gamma \sim 1$) $\tau c \sim 60$ km. Above the electron threshold, taking $m_{A^0} = 50$ MeV and $C_{Aff} = 1$, the total width is 10^{-5} eV and the corresponding decay length is $\tau c \sim 2$ cm.

Figure 1: $\tan\beta$ dependence of $\mathrm{BR}(A^0 \to e^+e^-)$. The lowest curve corresponds to $\tan\beta = 1$, the higher one to $\tan\beta = 3$, and the highest curve to $\tan\beta = 10$.

The $\tan\beta$ dependence of $\mathrm{BR}(A^0 \to e^+e^-)$ is shown in Fig. 1. By increasing $\tan\beta$ one reduces the up–type quark contributions to $\Gamma(A^0 \to \gamma\gamma)$, thereby increasing $\mathrm{BR}(A^0 \to e^+e^-)$. This dependence saturates for $\tan\beta \geq 3$. The decay mode $A^0 \to e^+e^-$ dominates for m_{A^0} below 80 MeV and switches off abruptly just above the electron threshold.

Below we study various constraints on $\{m_{A^0}, C_{Aff}\}$. An important class of constraints is due to the decays

$$X \to Y + \text{invisible}$$

with X, Y being some mesons. Experimental limits on their branching ratios exclude parts of parameter space where A^0 is sufficiently long-lived to escape the detector, that is $\tau\gamma c > d$ with τ being the lifetime of A^0, γ a boost factor and d the size of the detector (~ 10 m). In other words,

$$\Gamma_{\rm tot} \;<\; \frac{E_{A^0}}{m_{A^0}\, d}. \tag{15}$$

For a two-body decay $X \to Y + A^0$,

$$E_{A^0} = \frac{m_X^2 - m_Y^2 + m_{A^0}^2}{2m_X} \tag{16}$$

and in relevant cases $m_X^2 \gg m_Y^2, m_{A^0}^2$. Then for $m_{A^0} < 2m_e$, one has

$$m_{A^0}\sqrt{C_{Aff}} \;\lesssim\; 25\ \text{MeV} \times \sqrt[4]{\frac{m_X}{\text{GeV}}} \tag{17}$$

for $d \sim 10$ m and $\tan\beta \sim 1$. For $m_{A^0} > 2m_e$, one can estimate the resulting bound by taking $\Gamma_{\rm tot} \sim \Gamma_{e^+e^-}$, in which case

$$m_{A^0} C_{Aff} \;\lesssim\; 8\ \text{MeV} \times \sqrt{\frac{m_X}{\text{GeV}}}\ . \tag{18}$$

3.1 Rare B–decays $B \to K +$ invisible

Limits on production of A^0 in rare B-meson decays [31, 32, 33, 34] result from the following bounds on the branching ratios measured by CLEO [35] and BABAR [36]:

$$\begin{aligned} \mathcal{B}^{\rm CLEO}(B^0 \to K^0_S + \text{invisible}) &< 5.3 \times 10^{-5} \,, \\ \mathcal{B}^{\rm BABAR}(B^- \to K^- \nu\bar{\nu}) &< 7.0 \times 10^{-5} \,, \end{aligned} \tag{19}$$

where in what follows we use the more constraining CLEO result. Since in both experiments A^0 appears as missing energy, the bounds apply only if A^0 decays outside the detector. For $m_X = m_{B^0} = 5.28$ GeV, this implies

$$\begin{aligned} m_{A^0}\sqrt{C_{Aff}} &\lesssim 37 \quad \text{MeV} \quad \text{for } m_{A^0} < 2m_e \,, \\ m_{A^0}\, C_{Aff} &\lesssim 18 \quad \text{MeV} \quad \text{for } m_{A^0} > 2m_e \,. \end{aligned} \tag{20}$$

In the NMSSM, the decay rate for $B^0 \to K^0 A^0$ is given by [26]

$$\Gamma(B^0 \to K^0 A^0) = \frac{G_F^2\, |V_{tb}V_{ts}^*|^2}{2^{10}\pi^5}\, |C_A|^2\, \frac{|\vec{p}_K|}{m_{B^0}^2}\, \left| f_0^{B^0}(m_{A^0}^2) \right|^2 \left(\frac{m_{B^0}^2 - m_{K^0}^2}{m_b} \right)^2, \tag{21}$$

where the form factor $f_0^{B^0}(0) \sim 0.3 - 0.4$ [37] and $|\vec{p}_K| \simeq m_{B^0}/2$ is the three momentum of the kaon. The quantity C_A has been calculated in Ref. [26] in the large $\tan\beta$ regime, $C_A \sim C_{Aff} \tan\beta\; m_b m_t$ for order one stop mixing and EW scale sparticles. Since the full NMSSM calculation at low $\tan\beta$ is not available, we estimate the order of magnitude of the resulting bound by a rescaling of this result. Using the total B^0 width $\Gamma_{B^0} = 4.3 \times 10^{-13}$ GeV, the CLEO bound implies $C_{Aff} < 0.02/\tan\beta$. Taking conservatively $\tan\beta \sim \mathcal{O}(1)$, we get[1]

$$C_{Aff} < 10^{-2} \,. \tag{22}$$

This constraint is already strong at small $\tan\beta$ and gets even stronger at large $\tan\beta$.

The resulting exclusion region is shown in Fig. 2 (marked "$B^0 \to K^0 +$ inv."). We note that, in contrast to the lower boundary, the right boundary of this region is calculated quite reliably from Eq. 15 and is essentially independent of $\tan\beta$. In the plot, we use the full A^0-width $\Gamma_{\rm tot}$ without resorting to the approximation $\Gamma_{\rm tot} \sim \Gamma_{e^+e^-}$. The kink at $m_{A^0} \approx 2m_e$ is due to the rapid fall of $\Gamma(A^0 \to e^+e^-)$ as m_{A^0} approaches the threshold from above. Finally, the dependence on the detector size is only square–root.

3.2 Rare K–decays $K \to \pi +$ invisible

A light (invisible) A^0 can also be produced in K–decays. The relevant branching ratio has been measured by E787 [38, 39] and E949 [40]:[2]

$$\begin{aligned} \mathcal{B}^{\rm E787}(K^+ \to \pi^+ + \text{invisible}) &< 4.5 \times 10^{-11}, \\ \mathcal{B}^{\rm E949}(K^+ \to \pi^+ + \text{invisible}) &< 10^{-10}. \end{aligned} \tag{23}$$

[1] Essentially, this corresponds to the SM contribution with an additional coupling (9). The $b_R - s_L$ transition is mediated by the $W - t$ loop with A^0 coupled to the top quark.

[2] These bounds become significantly weaker at the pion pole, $m_{A^0} = m_\pi$.

We will use the tighter E787 bound. Eqs. 17 and 18 for $m_X = m_{K^+} = 494$ MeV yield the "invisibility" conditions

$$\begin{aligned} m_{A^0}\sqrt{C_{Aff}} &\lesssim 21 \quad \text{MeV} \quad \text{for } m_{A^0} < 2m_e \,, \\ m_{A^0}\, C_{Aff} &\lesssim 5 \quad \text{MeV} \quad \text{for } m_{A^0} > 2m_e \,. \end{aligned} \tag{24}$$

The decay rate is given by

$$\Gamma(K^+ \to \pi^+ A^0) = \frac{G_F^2\, |V_{ts}V_{td}^*|^2}{2^{10}\pi^5}\, |C_A'|^2\, \frac{|\vec{p}_\pi|}{m_{K^+}^2}\, \left| f_0^{K^+}(m_{A^0}^2) \right|^2 \left(\frac{m_{K^+}^2 - m_{\pi^+}^2}{m_s} \right)^2 , \tag{25}$$

where the form factor $f_0^{K^+}(0) \sim 1$ [41], $|\vec{p}_\pi| \simeq m_{K^+}/2$ and $C_A' \sim C_{Aff} \tan\beta\; m_s m_t$. With $\Gamma_{K^+} = 5.32 \times 10^{-17}$ GeV, the experimental bound requires $C_{Aff} < 2 \times 10^{-4}/\tan\beta$, which for $\tan\beta \sim \mathcal{O}(1)$ gives

$$C_{Aff} < 10^{-4} \,. \tag{26}$$

The corresponding excluded region is marked "$K^+ \to \pi^+ + \text{inv.}$" in Fig. 2. As in the case of B–decays, this bound only gets stronger with increasing $\tan\beta$ and its precise value is not important for us.

3.3 Rare decays $B \to K\; e^+e^-$, $K \to \pi\; e^+e^-$

If A^0 decays inside the detector, it contributes to the processes $B \to K\; e^+e^-$, $K \to \pi\; e^+e^-$. In experimental measurements, one sets a cut on the invariant mass $m_{e^+e^-} > 140$ MeV in order to suppress backgrounds including conversion photons and $\pi^0 \to e^+e^-\gamma$. Therefore, the resulting bounds apply for $m_{A^0} > 140$ MeV.

BELLE has reported [42]

$$\mathcal{B}^{\text{BELLE}}(B \to K\; \ell^+\ell^-) = (4.8^{+1.0}_{-0.9} \pm 0.3 \pm 0.1) \times 10^{-7} \,, \tag{27}$$

where ℓ includes muons and electrons with $m_{e^+e^-} > 140$ MeV. Assuming lepton universality, we will use a conservative bound $\text{BR}(B \xrightarrow{A^0} K\; e^+e^-) < 2.4 \times 10^{-7}$. The NMSSM result can be read off from Eq. 21 using

$$\text{BR}(B \xrightarrow{A^0} K\; e^+e^-) \simeq \text{BR}(B \to K\; A^0) \times \text{BR}(A^0 \to e^+e^-) \,. \tag{28}$$

In the relevant mass range, $\text{BR}(A^0 \to e^+e^-)$ is 20 - 40 % for $\tan\beta \sim 1$. After imposing the condition that A^0 decay inside the detector, we get

$$C_{Aff} < 8 \times 10^{-2} \,, \tag{29}$$

which excludes the strip 140 MeV $< m_{A^0} < 2m_\mu$ shown in Fig. 2 (marked "$B \to K e^+e^-$").

A similar result is obtained from K–decays [43]:

$$\mathcal{B}^{\text{NA48/2}}(K^\pm \to \pi^\pm\; e^+e^-) = (3.11 \pm 0.04 \pm 0.05 \pm 0.08 \pm 0.07) \times 10^{-7} \,, \tag{30}$$

Figure 2: Constraints from meson decays and muon $g-2$. The colored regions are excluded. These bounds include the effect of varying BR($A^0 \to e^+e^-$) with m_{A^0}.

which also employs the same kinematic cut $m_{e^+e^-} > 140$ MeV. Using Eq. 25 and an analog of (28), we get

$$C_{Aff} < 2 \times 10^{-2} \,, \tag{31}$$

where the "visibility" condition for A^0 has been imposed. This bound is shown in Fig. 2 marked "$K \to \pi e^+e^-$".

It is noteworthy that the window 140 MeV $< m_{A^0} < 2m_\mu$ is eliminated by 2 different processes. There are additional NMSSM contributions to $B \to K\ e^+e^-$, $K \to \pi\ e^+e^-$ apart from that of A^0, so in principle there could be cancellations. Considering two independent processes makes this possibility less likely.

The BaBar measurement of BR($B \to K\ \ell^+\ell^-$) [44] imposes a lower kinematic cut

$m_{e^+e^-} > 30$ MeV thereby losing somewhat in efficiency because of the low energy backgrounds [26]. Their result

$$\mathcal{B}^{\text{BABAR}}(B \to K\ \ell^+\ell^-) = (0.34 \pm 0.07 \pm 0.03) \times 10^{-6} \tag{32}$$

excludes the region $m_{A^0} > 30$ MeV with $C_{Aff} > 10^{-1} - 1$ (where A^0 decays inside the detector) depending on the pseudoscalar mass, Fig. 2 (marked "$B \to Ke^+e^-$"). Although one may question the reliability of this result at low e^+e^- invariant masses, another experiment, to be discussed in the next subsection, excludes a similar region of parameter space.

3.4 Rare K–decays $K^+ \to \pi^+ + X$

A byproduct of the $K_{\mu 2}$ experiment in Japan was a measurement of a 2-body decay $K^+ \to \pi^+ + X$, where X is any particle [45]. One searched for a peak in the π^+ momentum for 10 MeV $< m_X <$ 300 MeV. The resulting bound is

$$\mathcal{B}(K^+ \to \pi^+ + X) < 10^{-6} \tag{33}$$

at 90% CL for $m_X < 60$ MeV. For larger m_X up to 120 MeV this bound relaxes to 10^{-5}.

The excluded parameter space is shown in Fig. 2 (marked "$K^+ \to \pi^+ X$"). The constraint amounts approximately to

$$C_{Aff} < 4 \times 10^{-2} \tag{34}$$

for $m_{A^0} > 10$ MeV.

We note that at $m_{A^0} \simeq m_{\pi^0}$ the constraint is weaker. However, this region is disfavored by $\pi^0 \to e^+e^-$ (see below) and the $\pi^+ - \pi^0$ mass difference when $\pi^0 - A^0$ mixing is taken into account.

A similar region of parameter space (up to $m_{A^0} = 100$ MeV) is excluded by the process $\pi^+ \to e^+\nu\ A^0$ with subsequent decay $A^0 \to e^+e^-$ [46] (for applications to axion models, see [47]).

3.5 Radiative Upsilon–decays

The bounds on A^0 production in radiative Υ decays come from CLEO [48] and BABAR [49, 50]:

$$\mathcal{B}^{\text{CLEO}}(\Upsilon(1S) \to \gamma + \text{invisible}) \ < \ 1.3 \times 10^{-5}, \tag{35}$$

$$\mathcal{B}^{\text{BABAR}}(\Upsilon(3S) \to \gamma + \text{invisible}) \ < \ 3 \times 10^{-6}. \tag{36}$$

They apply only if A^0 decays outside the detector, which for $m_X = m_{\Upsilon(3S)} = 10.4$ GeV means

$$\begin{aligned} m_{A^0}\sqrt{C_{Aff}} &\lesssim 44 \ \text{MeV} \quad \text{for } m_{A^0} < 2m_e\ , \\ m_{A^0}\ C_{Aff} &\lesssim 25 \ \text{MeV} \quad \text{for } m_{A^0} > 2m_e\ . \end{aligned} \tag{37}$$

The branching ratio for $\Upsilon \to A^0\gamma$ is given by [51, 52, 26]

$$\frac{\mathcal{B}(\Upsilon \to A^0\gamma)}{\mathcal{B}(\Upsilon \to \mu^+\mu^-)} = \frac{G_F m_b^2}{\sqrt{2}\pi\alpha}\ C_{Aff}^2\ \left(1 - \frac{m_{A^0}^2}{m_\Upsilon^2}\right)\ F_{\text{QCD}}\ , \tag{38}$$

where $F_{\rm QCD} \sim 0.5$ is a QCD correction factor, $\mathcal{B}(\Upsilon(1S) \to \mu^+\mu^-) = 0.025$ and $\mathcal{B}(\Upsilon(3S) \to \mu^+\mu^-) = 0.022$ [53]. The resulting bounds (marked in Fig. 2 "$\Upsilon(1S) \to \gamma + \text{inv.}$" and "$\Upsilon(3S) \to \gamma + \text{inv.}$") are

$$\begin{aligned} C_{Aff} &< 0.37 \quad (\text{CLEO}) \,, \\ C_{Aff} &< 0.19 \quad (\text{BABAR}) \,, \end{aligned} \tag{39}$$

independent of $\tan\beta$.

3.6 Pion decay $\pi^0 \to e^+e^-$

A light A^0 provides a pseudoscalar channel for pion annihilation into e^+e^- (see, e.g. [54]). This chirality–suppressed decay proceeds in the SM through a loop diagram with a $\pi^0\gamma\gamma^*$ vertex and has a very small branching ratio. The recent KTEV result [55]

$$\mathcal{B}^{\rm KTeV}(\pi^0 \to e^+e^-) = (7.48 \pm 0.29 \pm 0.25) \times 10^{-8} \tag{40}$$

is somewhat (3σ) above the SM prediction [56, 57]. To be conservative, in what follows we will require that the tree level contribution from A^0 not exceed the central experimental value, $\mathcal{B}(\pi^0 \xrightarrow{A^0} e^+e^-) < 7.5 \times 10^{-8}$.

We find

$$\Gamma(\pi^0 \xrightarrow{A^0} e^+e^-) \;\simeq\; \frac{G_F^2}{4\pi} \, \frac{m_e^2 m_\pi^5 f_\pi^2}{|m_\pi^2 - m_{A^0}^2 + i\Gamma_{A^0} m_{A^0}|^2} \, C_{Aff}^4 \,, \tag{41}$$

where we have neglected the up–quark contribution, $\langle 0|m_d \bar{d}\gamma^5 d|\pi^0\rangle \simeq -i m_\pi^2 f_\pi$.[3] Γ_{A^0} is given by Eq. 10, $m_\pi = 135$ MeV and $f_\pi = 93$ MeV. We neglect the $\pi^0 - A^0$ mixing effects which are of order $\delta m^2/m_\pi^2 \sim f_\pi/M_W \sim 10^{-3}$ and relevant only very close to the pion mass.

The total width of π^0 is

$$\Gamma(\pi^0 \to \gamma\gamma) = \frac{\alpha^2}{64\pi^3} \, \frac{m_\pi^3}{f_\pi^2} \,, \tag{42}$$

then the KTEV result requires

$$C_{Aff} < 20 \tag{43}$$

away from the resonance region, where the constraint is stronger. A more precise bound including the m_{A^0}–dependence is shown in Fig. 2 marked "$\pi^0 \to e^+e^-$". This constraint is complementary to those of the $X \to Y +$ invisible decays, in that it excludes parameter space above $C_{Aff} \simeq 20$ regardless of the A^0 mass. It is also a reliable tree–level constraint essentially independent of $\tan\beta$.

[3] Our bound on C_{Aff} is not sensitive to this approximation as it scales as a square root of this matrix element.

3.7 Muon anomalous magnetic moment

At loop level, A^0 contributes to the muon $g-2$ which is well measured. Currently, there is a 4σ discrepancy between the SM prediction for the muon $g-2$ and its measured value at BNL E821 [58]:

$$\Delta a_\mu = a_\mu^{\rm Exp} - a_\mu^{\rm SM} = (31.6 \pm 7.9) \times 10^{-10} \,, \tag{44}$$

which may be considered a hint for new physics.

In the NMSSM, there are significant one- and two-loop contributions of the CP-odd Higgs A^0 to a_μ. They are given, for example, in [59]:

$$\begin{aligned}
\delta a_\mu(A^0) &= \delta a_\mu^{\rm 1L}(A^0) + \delta a_\mu^{\rm 2L}(A^0) \,, \\
\delta a_\mu^{\rm 1L}(A^0) &= -\frac{\sqrt{2} G_F}{8\pi^2}\, m_\mu^2\, |C_{Aff}|^2\, f_1\Big(\frac{m_{A^0}^2}{m_\mu^2}\Big) \,, \\
\delta a_\mu^{\rm 2L}(A^0) &= \frac{\sqrt{2} G_F \alpha}{8\pi^3}\, m_\mu^2\, |C_{Aff}|^2 \left[\frac{4}{3}\frac{1}{\tan^2\beta} f_2\Big(\frac{m_t^2}{m_{A^0}^2}\Big) + \frac{1}{3} f_2\Big(\frac{m_b^2}{m_{A^0}^2}\Big) + f_2\Big(\frac{m_\tau^2}{m_{A^0}^2}\Big)\right] \,,
\end{aligned} \tag{45}$$

where

$$\begin{aligned}
f_1(z) &= \int_0^1 dx \frac{x^3}{x^2 + z(1-x)} \,, \\
f_2(z) &= z \int_0^1 dx \frac{1}{x(1-x) - z} \ln\frac{x(1-x)}{z} \,.
\end{aligned} \tag{46}$$

The one loop-contribution is negative which makes the discrepancy worse. For m_{A^0} above roughly 1 GeV, the two–loop contribution may be dominant and resolve the discrepancy. However, this does not occur in the mass range we consider.

Since there are NMSSM contributions to $g-2$ of both signs, the contribution from the CP–odd Higgs A^0 can be canceled. We then require that the latter not worsen the discrepancy beyond 5 σ:

$$\delta a_\mu(A^0) \leq a_\mu^{\rm Exp} - a_\mu^{\rm SM} \simeq 40 \times 10^{-10} \quad (5\sigma) \,. \tag{47}$$

The corresponding bound on C_{Aff} (at $\tan\beta \sim 1$) is shown in Fig. 2 marked "a_μ". It can be approximated by

$$C_{Aff} < 2 \tag{48}$$

for $m_{A^0} \lesssim m_\mu$. The $\tan\beta$ dependence is very mild in the region of interest and stems only from the 2–loop contribution, which is subdominant.

Once this bound is imposed, the electron $g-2$ constraint is satisfied automatically.

3.8 Other constraints

Further (model–dependent) constraints are summarized in Refs. [24, 26]. These are weaker than the bounds we have considered and require assumptions about the NMSSM spectrum. For instance, there are contributions from all neutral Higgses to $B_s \to \mu^+\mu^-$ and $B - \bar{B}$

mixing which allows one to eliminate parts of parameter space ($C_{Aff} \sim \mathcal{O}(10)$) depending on their masses and $\tan\beta$ [26].

Among other possible flavor physics constraints are J/Ψ decays. CLEO has recently reported $\mathcal{B}^{\rm CLEO}(J/\Psi \to \gamma + \text{invisible}) < 4.3 \times 10^{-6}$ [60], which is somewhat weaker than the analogous $\Upsilon(3S)$ bound. The A^0 coupling to up–type quarks falls very quickly with $\tan\beta$, so we do not use this result in our analysis.

Further, the missing–energy process $B \to K A^0 A^0$ proceeding through the $h A^0 A^0$ coupling [61] sets a mild constraint on the $S H_1 H_2$ coupling in the superpotential, $\lambda < 0.7$.

A light CP–odd Higgs could potentially be constrained by the LEP data. A^0 couples to the Z–boson at tree level through the $A^0 H_i^0 Z_\mu$ vertex [24], therefore the (invisible) Z–width does not constrain the mass of A^0. Furthermore, the electroweak oblique corrections are suppressed by the mass of the heavier pseudoscalar (see, e.g. [15]). For the same reason, the A^0 production at LEP through $e^+e^- \to h\ A^0$ is suppressed. The associated production with bottom quarks $e^+e^- \to b\bar{b}\ A^0$ is also insignificant [15]. Finally, the constraints from $Z \to \gamma\ A^0$ are weak [62, 63].

Astrophysical bounds have been summarized in Ref. [20]. They are usually relevant for sub–MeV pseudoscalar masses which, in the range $10^{-4} < C_{Aff} < 10^3$, are already excluded by meson decays with missing energy. However, the supernova SN1987A sets an additional constraint for a small coupling: $C_{Aff} > 10^{-4}$ when $m_{A^0} < 30$ MeV [20].

To summarize this section, we see that a combination of various constraints requires the CP–odd Higgs to be heavier than $2m_\mu$ (unless $C_{Aff} < 10^{-4}$). To obtain this bound we did not rely on the specifics of the NMSSM. All we used was the coupling (9) at $\tan\beta \sim \mathcal{O}(1)$, which is much more general. This coupling is sufficient to induce the $b-s$ and $s-d$ transitions (with flavor change due to SM loops) which were used in the processes like $B \to K\ A^0$ and $K \to \pi\ A^0$. Similarly, Υ decays, $\pi^0 \to e^+e^-$ and muon $g-2$ are generated directly by (9).

For completeness, in the next section we discuss the reactor and beam dump results which have been used in the past to constrain axion models.

4 Further bounds from reactor and beam dump experiments

4.1 Reactor bounds

Searches for axion–like particles using nuclear power reactors set constraints on the parameter space of the CP–odd Higgs. Here we consider 2 representative experiments which employ Bugey and Kuo–Sheng nuclear reactors.

Axion–like particles can be emitted in place of photons from excited nuclear levels which makes nuclear reactors a source of pseudoscalars with masses up to 10 MeV. In [64], the detector was placed 18.5 m from the Bugey reactor core and one searched for the decays $A^0 \to e^+e^-$. No excess of e^+e^- events has been observed which set a constraint on the axion decay constant f_χ. The corresponding exclusion region can be read off from Fig. 5 of [64] using the conversion

$$C_{Aff} = \frac{1}{f_\chi}\ \frac{2m_W}{g}\ . \tag{49}$$

The result is shown in Fig. 3.

Figure 3: Constraints from nuclear power reactors.

Another experiment at the Kuo–Sheng nuclear reactor searched for axions via Compton conversion on electrons [65]. A Germanium detector, placed 28 m away from the reactor, measured the ionization energy resulting from the axion–photon conversion in the detector. Again, no signal was found. The exclusion region can be read off from Fig. 7 of [65] and their Eq. 31, $g_{aee}\, g^1_{aNN} < 1.3\times10^{-10}$. Using $g_{aee} = C_{Aff} g m_e/(2m_W)$ and $g^1_{aNN} = 3\times10^{-8} m_{A^0}/\mathrm{eV}$, this translates into

$$m_{A^0}\; C_{Aff} < 2 \times 10^{-3}\ \mathrm{MeV} \tag{50}$$

for $m_{A^0} < 2m_e$. The experiment is also sensitive to m_{A^0} up to 2.23 MeV, which is the $pn \to d\gamma$ transition energy. Requiring that the axion not decay before it reaches the detector ($m_{A^0} C_{Aff} < 0.3$ MeV), one obtains the bulge at $m_{A^0} > 2m_e$ in Fig. 3.

4.2 Beam dump limits

Axion–like particles can be emitted via bremsstrahlung or Primakoff production in beam dump experiments (see, e.g. [66]). The setup of these experiments is as follows. An intense beam of particles (electrons or protons) hits a thick target, which absorbs the beam and the interaction products apart from very weakly interacting particles such as axions. The decay products of the latter are collected by the detector, typically placed tens of meters behind the target.

Below we consider 4 representative beam dump experiments.[4]

- SLAC E141 [68]: 2×10^{15} electrons at energy 9 GeV struck a 12 cm tungsten target; detector 35 m behind the target
- Fermilab E774 [69]: 0.52×10^{10} electrons at 275 GeV dumped at a 30 cm target; detector length 7.25 m
- CHARM [70]: 2.4×10^{18} protons at 400 GeV dumped at a thick copper target; detector 480 m behind the target
- Orsay [71]: 2×10^{16} electrons at energy 1.6 GeV dumped in a 1 m target; detector 2 m behind the target

The corresponding exclusion regions (Fig. 4) can be read off from the plots presented in these papers either using the conversion factor for the axion decay constant (49) or calculating the axion decay time according to Eq. 10.

Figure 4: Constraints from the beam dump experiments.

The reactor and beam dump experiments by themselves eliminate most of the parameter space (Figs. 3,4). They are based on a different kind of physics compared to meson decays

[4]We are not displaying the results of the SLAC E137 [66] and Fermilab 605 [67] experiments since the corresponding exclusion regions are largely covered by other experiments.

of Sec. 3 and in this sense are complementary. Note also that, unlike meson decays, electron beam dump experiments as well as the muon $g-2$ probe directly the lepton–axion coupling, which could be the only axion coupling to matter in exotic ("leptophilic") scenarios.

5 Conclusion

In this work, we have studied the question how light a CP–odd Higgs A^0 of the NMSSM can be. We have analyzed constraints from meson decays, muon $g-2$, beam dump and reactor experiments. We find that the parameter space $m_{A^0} < 2m_\mu$ is excluded (by more than one experiment) unless the coupling of A^0 to matter is 4 orders of magnitude smaller than that of the Standard Model Higgs, i.e. $C_{Aff} < 10^{-4}$. Since such a small coupling can hardly be achieved in the NMSSM, we conclude that A^0 has to be heavier than about 210 MeV.

Our analysis applies more generally to couplings of a light pseudoscalar to matter. We have not used any specific features of the NMSSM (nor supersymmetry). We have only relied on Eq. 9 and analyzed parameter space in terms of $\{m_{A^0}, C_{Aff}\}$. Flavor changing couplings of the Standard Model are sufficient to generate the processes like $B \to K\ A^0$ and $K \to \pi\ A^0$, which lead to strong constraints.

Since the CP–odd Higgs is heavier than $2m_\mu$, the decay channel to muons is open. One can therefore produce A^0 in gluon fusion at the LHC and look for $\mu^+\mu^-$ pairs with low invariant mass [72], or search for decays $h \to 2A^0 \to 4\mu$ [73, 74] (similar to what has been done by D0 at Tevatron [75]). This will provide an important test of models with a light pseudoscalar.

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