Higgs Boson Decays to CP-odd Scalars at the Tevatron and Beyond

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In extended Higgs models, the Higgs boson may decay into a pair of light CP-odd scalars, with distinctive collider signatures. We study the ensuing Higgs signals at the upgraded Tevatron, considering the subsequent decays of the scalars into pairs of gluons or photons. For CP-odd scalars lighter than a few GeV, the Higgs boson manifests itself as a diphoton resonance and can be discovered up to masses of a few hundred GeV. For heavier CP-odd scalars the reach extends at most up to $M_h \sim 120$ GeV. We also discuss the capabilities of the LHC and lepton colliders in these channels.

The Higgs boson is the only Standard Model (SM) particle that remains elusive. The next runs at the Fermilab Tevatron have the potential for discovering the Higgs boson for a mass range beyond the current LEP limit [1]. Remarkable efforts [2] have been devoted to designing the best search strategy for a SM-like Higgs boson, with a large branching into b-jets or W-bosons. However, a Higgs boson, h^0, with SM-like couplings to the gauge bosons and fermions, could nevertheless have decay modes dramatically different from the SM ones. The reason is that the SM is likely to be a part of a more comprehensive theory which may include an extended Higgs sector. It is then possible for the Higgs boson to decay into pairs of other neutral scalars, whenever they are lighter than half the h^0 mass, M_h.

An example of such an extended Higgs sector is given by the Minimal Composite Higgs Model (MCHM) [3], which is based on the top-quark condensation seesaw mechanism [4,5]. At low energy, the MCHM includes two composite Higgs doublets and two gauge-singlet scalars, with the h^0 and a CP-odd scalar, A^0, being the lightest scalar mass eigenstates. Another interesting example is the Next-to-Minimal Supersymmetric Standard Model (NMSSM) [6], where the presence of two Higgs doublets and a gauge singlet allows a region of parameter space in which M_h is larger than twice the mass M_A of the lightest CP-odd scalar, A^0. Both in the MCHM and NMSSM it is natural to have a light A^0 because its mass is controlled by the explicit breaking of a spontaneously broken $U(1)$ symmetry. In both models this global $U(1)$ symmetry has a QCD anomaly, and therefore A^0 is an axion. Note though that various axion searches [7] place a lower bound on M_A, typically above ~ 1 MeV, which requires explicit $U(1)$ breaking beyond the QCD anomaly, so that A^0 does not solve the strong CP problem. In other models, such as the chiral supersymmetric models [8], or composite Higgs models from extra dimensions [9], h^0 could also decay into light CP-even scalars.

In this Letter we study the Higgs boson decay into CP-odd scalar pairs at the upgraded Tevatron. We assume the existence of a scalar A^0 (we call it "axion" for short),

of mass $M_A < M_h/2$, with a trilinear coupling

$$\frac{c\,v}{2} h^0 A^0 A^0 \,, \tag{1}$$

where $v \approx 246$ GeV is the electroweak symmetry breaking scale, and c is a model-dependent dimensionless parameter, which can be as large as $\mathcal{O}(1)$.

The Higgs width into a pair of axions is

$$\Gamma(h^0 \to A^0 A^0) \;=\; \frac{c^2\,v^2}{32\pi M_h}\left(1 - 4\frac{M_A^2}{M_h^2}\right)^{1/2}. \tag{2}$$

The decay to axion pairs can be essential for Higgs boson searches in collider experiments. This is illustrated in Fig. 1, where we plot the Higgs boson branching ratio to axions, $B(h^0 \to A^0 A^0)$, versus M_h, for $M_A \ll M_h/2$ and several values of c. For M_h below the WW threshold, the dominant SM decay of the Higgs boson $h^0 \to b\bar{b}$ has a very small width, and is therefore susceptible to the presence of new physics beyond the SM, *e.g.* the interaction (1). The decay to axions would then dominate over $h \to b\bar{b}$ for values of c as small as $\sim 0.02\,(M_h/100\text{ GeV})$. We also see that even above the WW threshold, $h^0 \to A^0 A^0$ competes with $h^0 \to WW$, provided $c \sim \mathcal{O}(1)$.
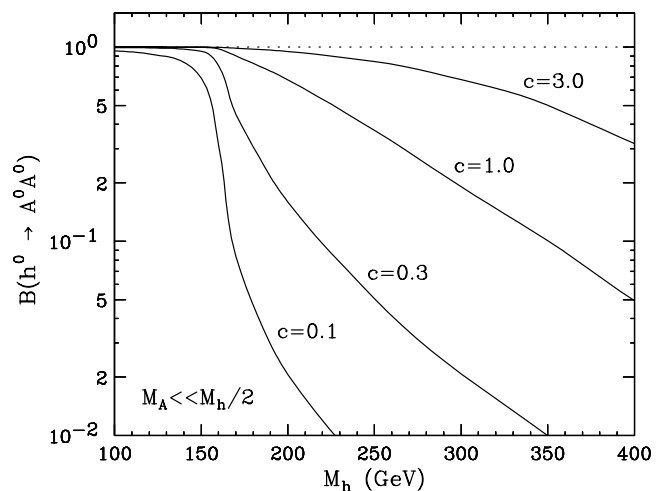

FIG. 1. Branching ratio of the Higgs boson into axion pairs, as a function of M_h, for $M_A \ll M_h/2$ and several values of c.

Experimental signatures of the Higgs boson.— The final state signatures for $h^0 \to A^0 A^0$ searches will depend on the subsequent decays of the axions, which are quite model-dependent. Here we will concentrate on the case where the axion couplings to light fermions are negligible. Such a situation may arise in the MCHM, *e.g.* when only one Higgs doublet is mainly responsible for electroweak symmetry breaking and fermion masses (type I two-Higgs-doublet model with large $\tan\beta$ [6]). In the MCHM, however, the axion has a large coupling to a heavy vector-like quark, χ, whose charges under the SM gauge group are the same as for the right-handed top quark t_R. This coupling allows one-loop decays of the axion into gluon or photon pairs. The case where the axion couplings to light fermions are significant (*e.g.*, in the NMSSM) provides interesting experimental signatures as well, but we leave their investigation for future studies.

The constraints on a light CP-odd scalar with small couplings to SM fermions are loose [3]. For example, no constraints on the axion mass have been set from $Z \to A^0 \gamma$ at LEP [10], from direct A^0 production through gluon fusion via a χ loop at the Tevatron [11], from fits to the electroweak data, or from meson decays [6]. The relevant lower bounds on M_A come from beam dump experiments [12], in the MeV range, and from star cooling rates, $M_A \gtrsim 0.2$ MeV [7].

In the MCHM, the two Higgs doublets and two gauge-singlet scalars arise as bound states of the top-bottom left-handed doublet or χ_L, with χ_R or t_R. At scales below a few TeV this is an explicit, renormalizable theory with scalars that *appear* to be fundamental. The Higgs potential has been analyzed in detail in [3,5]. For the purpose of studying Higgs boson decays in the MCHM it is however sufficient to introduce the following simpler model. Consider the SM with the addition of a gauge singlet complex scalar S and the vector-like quark, χ. Besides the kinetic terms and Higgs Yukawa couplings, the following terms are present in the Lagrangian:

$$\mathcal{L} = \xi S \overline{\chi}_L \chi_R + \text{h.c.} - V(H, S) \ , \quad (3)$$

where ξ is a Yukawa coupling, and the scalar potential is

$$V(H,S) = \frac{\lambda_H}{2} \left(H^\dagger H\right)^2 + \frac{\lambda_S}{2} \left(S^\dagger S\right)^2 + \lambda_0 H^\dagger H S^\dagger S$$
$$+ M_H^2 H^\dagger H + M_S^2 S^\dagger S + C_S \left(S + S^\dagger\right) \ . \quad (4)$$

We assume $M_H^2 < 0$, and $\lambda_H \lambda_S > \lambda_0^2$. In the limit where the coefficient C_S of the tadpole term vanishes, the effective potential has a global $U(1)$ symmetry spontaneously broken by the vacuum expectation value $\langle S \rangle$ of S. The associated axion, A^0, is part of the singlet S, and due to the tadpole term has a mass given by $\sqrt{|C_S/\langle S \rangle|}$. Note that in the MCHM the tadpole term is generated by a tree level χ mass, and there is mixing between χ and t once the electroweak symmetry is broken. These details are not relevant for the present study, but the coupling of

the axion to χ is essential since otherwise the axion would be stable and the Higgs boson would decay invisibly [13].

The $h^0 A^0 A^0$ coupling may be easily computed in the SM+singlet model, with the result

$$c = -\sqrt{2} \lambda_S \frac{\langle S \rangle}{v} \sin\theta - \lambda_0 \cos\theta \ , \quad (5)$$

where θ is the mixing angle between the two CP-even neutral scalars,

$$\tan 2\theta \simeq -\frac{2\sqrt{2} \lambda_0 v \langle S \rangle}{2\lambda_S \langle S \rangle^2 - \lambda_H v^2 - C_S/\langle S \rangle} \ . \quad (6)$$

For a range of values of the six parameters from Eq. (4), the Higgs decay to axions is important.

The effective axion couplings to gluons and photons, induced at one-loop by the χ quark, are given by

$$\frac{\sqrt{2}}{16\pi \langle S \rangle} A^0 \epsilon^{\mu\nu\rho\sigma} \left(\alpha_s \mathbf{G}_{\mu\nu} \mathbf{G}_{\rho\sigma} + N_c e_\chi^2 \alpha F_{\mu\nu} F_{\rho\sigma} \right) \ , \quad (7)$$

where α_s (α) is the strong (fine structure) constant, $\mathbf{G}_{\mu\nu}$ ($F_{\mu\nu}$) is the gluon (photon) field strength, $N_c = 3$ is the number of colors and $e_\chi = 2/3$ is the electric charge of the vector-like quark. Naively, the dominant axion decay mode is to a pair of gluons. However, the gluons must then hadronize, and for light axions the number of open channels is very limited. Indeed, the two-body decays $A^0 \to \pi\pi, \pi^0 \gamma$ are forbidden by CP invariance and angular momentum conservation, while the three body decays $A^0 \to \pi^0 \gamma\gamma, 2\pi^0 \gamma, \pi^+ \pi^- \gamma$ are significantly suppressed by phase space (note that C- and P-parity need not be separately conserved in axion decays). Therefore, for $M_A \lesssim 3m_\pi \approx 405$ MeV the branching ratio to two photons is close to 1. The axion decay width due to the two photon effective coupling (7) is given by

$$\Gamma(A^0 \to \gamma\gamma) \simeq \frac{\alpha^2 M_A^3}{72\pi^3 \langle S \rangle^2} \ . \quad (8)$$

For $M_A \gtrsim 0.5$ GeV the isospin-violating decay modes $A^0 \to 3\pi$ open up and start competing with $A^0 \to \gamma\gamma$. Due to QCD uncertainties, it is quite difficult to estimate their exact width, but recall that the measured branching fractions of η decays into $\gamma\gamma$, $3\pi^0$, and $\pi^+\pi^-\pi^0$ are given by 39%, 32%, and 23%, respectively [7]. Since η and A^0 have the same quantum numbers, the $A^0 \to \gamma\gamma$ decay mode is likely to be significant even for M_A of order 1 GeV. Fortunately, the study of the $h^0 \to A^0 A^0$ decay in this M_A range does not require a precise determination of the axion branching ratios, as we explain below.

The Higgs decays are quite peculiar in the scenario discussed here. Since the LEP [14] and Tevatron [15] limits would generally apply, we only consider the range of Higgs masses above ~ 100 GeV. Because of the relatively heavy parent mass, each axion will be produced with a significant boost, and will decay into a pair of

almost colinear photons. For $M_A \lesssim 0.025 M_h$, they will not be resolved in the electromagnetic calorimeter and will be reconstructed as a single photon. As a result, the $h \to A^0 A^0 \to 4\gamma$ mode will appear in the detector as a diphoton signature, as shown in Fig. 2a.
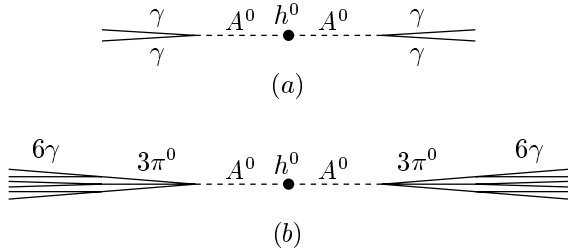
FIG. 2. Higgs boson decay topology into a "diphoton" final state with (a) prompt and (b) cascade photons.

The interesting twist is that the $A^0 \to 3\pi^0$ decay mode will have the same signature in the detector, because π^0's decay promptly into photons. This situation is illustrated in Fig. 2b. We can continue this line of argument even for higher M_A. If $M_A \gtrsim 1$ GeV, the $\omega\gamma$, $\rho\gamma$, and $\eta\pi\pi$ axion decay modes become relevant. Using the measured branching fractions of η' into these states, and the subsequent η, ρ, ω decays, we find that the η' branching fraction for final states with only photons is roughly 17%. Given the similarities between η' and the axion, we expect the probability that A^0 is reconstructed as a photon is of order 20% when $M_A \sim 1$ GeV.

For $M_A \gtrsim 2$ GeV, new isospin-conserving modes with large branching fractions open up: $A^0 \to \rho\rho$, $\omega\omega$, $a_0\pi^0$, $KK\pi^0$, etc. (isospin-violating decays, such as $\eta\eta\pi$, are suppressed.) Even then, some of these mesons have large branching fractions into states which subsequently decay into photons, e.g., $a_0 \to \eta\pi^0$ [7], yielding the same diphoton signature for h^0. Only when M_A is increased above several GeV, the decay products of h^0 start looking more and more like QCD jets instead of photons.

An important issue is whether the A^0 decays promptly. Using Eq. (8), we can estimate its decay length

$$L_A \approx 4 \text{ mm} \frac{M_h \langle S \rangle^2}{(100 \text{ GeV})^3} \left(\frac{1 \text{ GeV}}{M_A} \right)^4 . \qquad (9)$$

For $M_A \lesssim 100$ MeV the axion decays occur most of the time outside the detector, and the Higgs boson decay is invisible [16]. In the following we shall instead concentrate on the mass range $M_A \gtrsim 200$ MeV, where the axion decays before reaching the electromagnetic calorimeter.

Tevatron reach for light axions. — For M_A less than a few GeV, we have studied the Tevatron reach in the diphoton channel, following the analysis of Ref. [17]. We used PYTHIA [18] for event generation and SHW [19] with minor modifications [20] for detector simulation. The Higgs boson is produced predominantly via gluon

fusion and the inclusive diphoton channel (with an optimized cut on the invariant diphoton mass, $m_{\gamma\gamma}$) provides the best sensitivity [17]. The reach is shown in Fig. 3, where we plot the product $L \times P^2(A \to \gamma)$ as a function of M_h. Here L is the total integrated luminosity required for a 95% confidence level exclusion, and $P(A \to \gamma)$ is the (M_A-dependent) probability that an axion is reconstructed as a photon, including the branching fractions of cascade decays of axions into multi-photon final states. We see that already run IIa with 2 fb^{-1} will be probing a range of Higgs boson masses well beyond the reach via conventional SM searches [2]. Note the reversed ordering of the lines of constant c at low and high M_h. Below the WW-threshold, the Higgs boson width, Γ_h, is dominated by the decay (2) into axions, hence a larger c requires a softer $m_{\gamma\gamma}$ cut and leads to higher background. Above the WW-threshold, where Γ_h is dominated by the SM decays, a larger c is beneficial due to the increased $B(h \to A^0 A^0)$ (see Fig. 1).
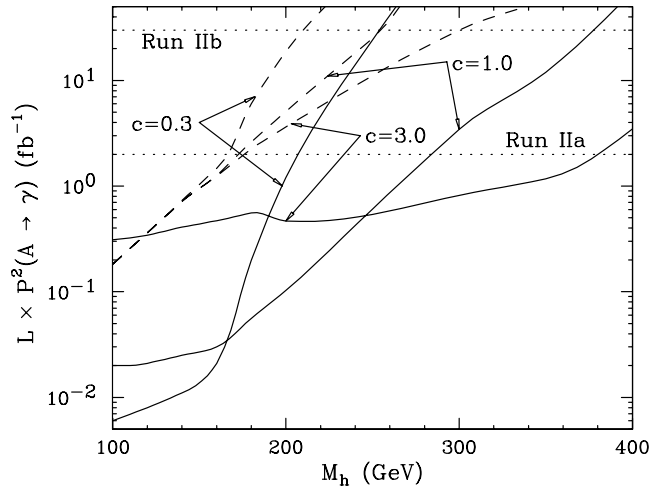
FIG. 3. The Tevatron reach in the inclusive diphoton channel (solid) and $\ell 2\gamma \not{E}_T$ channel (dashed), as a function of M_h, for three different values of c.

For light axions, associated Wh/Zh production can give alternative, very clean signatures: $\ell\gamma\gamma\not{E}_T$ and $\ell^+\ell^-\gamma\gamma$. Considering only leptonic decays of the W/Z, we expect Wh to give a better reach, because of the larger leptonic branching fraction. Requiring events with $\not{E}_T > 20$ GeV, at least one lepton with $p_T(\ell) > 20$ GeV, and at least two photons with $p_T(\gamma) > 20$ GeV and $|m_{\gamma\gamma} - M_h| < \Gamma_h$, we find the following parametrization of the signal efficiency: $\varepsilon = 0.32 - 0.07(M_h/100 \text{ GeV})$. The main background is $W\gamma\gamma$ and we estimated it using COMPHEP [21] to be less than 1 event after cuts in 30 fb^{-1}. The corresponding reach is shown in Fig. 3 (dashed lines). We see that although this channel gives a smaller absolute reach than the inclusive 2γ channel, it will still enable the Tevatron to probe Higgs masses below WW threshold, for a wide range of values of c.

Tevatron reach for heavy axions. — If M_A is above a

few GeV, the decay products from $A \to gg$ look like QCD jets, and the background is large. To make matters worse, at the parton level $A^0 \to gg$ dominates over $A^0 \to \gamma\gamma$ by a factor of $(N_c^2 - 1)\alpha_s^2/(4N_c^2\alpha^2 e_\chi^4) \approx 265$. Typically, for $M_A \gtrsim 20$ GeV, the resulting two gluon jets are separated well enough and can both be individually reconstructed. Unlike the SM, the final states from Higgs production contain no b-jets, but instead mostly gluon jets and very rarely photons. Since at present it is practically impossible to distinguish between gluon and quark jets, these searches appear quite challenging.

Gluon fusion ($gg \to h \to A^0 A^0$) can lead to $4j$, $2j2\gamma$ or 4γ events. The $4j$ channel suffers from insurmountable backgrounds, while the 4γ channel has a tiny branching ratio, and will escape detection even in run IIb with 30 fb^{-1}. The $2j2\gamma$ channel is similar to the searches for a "bosonic Higgs" at the Tevatron [15,17,22]. Using the same cuts as for the inclusive $\gamma\gamma + N$jets channels of Ref. [17], we found no reach in run IIb.

Searches for associate Higgs boson production also appear very difficult. The 6-jet final states from all-hadronic decays of the W (Z) and A^0 have large QCD backgrounds. Leptonic decays of the W and Z, combined with $h \to 4j$ give $\ell 4j\!\!\!/E_T$ and $\ell^+\ell^- 4j$, respectively. The backgrounds are large and again no sensitivity in run II is expected. Finally, requiring that one axion decays to $\gamma\gamma$, and leptonic decays of the W or Z, we get the relatively clean $\ell 2\gamma\!\!\!/E_T + X$ and $\ell^+\ell^- 2\gamma + X$ final states. Because of the $B(A^0 \to \gamma\gamma)$ suppression, the signal rates are too small to be observed in run IIa, but run IIb might be able to explore the mass range $100 \lesssim M_h \lesssim 120$ GeV.

Discovery prospects at the LHC and future lepton colliders.— It is interesting to contemplate the capabilities of future colliders for our scenario. The LHC has enormous potential for such Higgs boson searches. Just like at the Tevatron, one will have to concentrate on the cleanest channels. But the much larger signal rate will allow to look for the photonic decays of heavy axions, which were severely limited by statistics at the Tevatron. Preliminary studies show that even the $2j2\gamma$ channel, swamped by the QCD background at the Tevatron, is sensitive to a large range of Higgs boson and axion masses at the LHC, due to the improved energy resolution of the LHC detectors and the enhanced signal cross section.

A high energy lepton collider (such as the NLC) would be ideally suited for unravelling a non-standard Higgs sector, like the one discussed in this Letter, particularly if M_A is bigger than a few GeV and $P(A \to \gamma)$ is very small. Notice that the jet-rich channels are the best to look for at a lepton collider, since they have the largest branching ratios. In this sense, lepton colliders are complementary to hadron machines, where these channels suffer from large backgrounds. We also expect that LEP-II will be able to probe Higgs masses up to its kinematic reach, once a dedicated search is done. We therefore urge the LEP collaborations to present Higgs mass limits with the data selection optimized for the discussed signatures.

In conclusion, we have considered several novel Higgs boson discovery signatures, arising from the decay $h^0 \to A^0 A^0$, present in many extended Higgs sector models. Quite ironically, the best reach at the Tevatron is obtained for light axions, where the Higgs boson often decays to two jets, each of which "fakes" a photon.

[1] E. Ferrer, talk given at "XXXV Rencontres de Moriond", Les Arcs, March 11-18 2000.

[2] Report of the Higgs Working Group, conveners M. Carena, J. Conway, H. Haber, and J. Hobbs, http://fnth37.fnal.gov/higgs.html.

[3] B. A. Dobrescu, hep-ph/9908391.

[4] B. A. Dobrescu and C. T. Hill, Phys. Rev. Lett. **81**, 2634 (1998).

[5] R. S. Chivukula, B. A. Dobrescu, H. Georgi and C. T. Hill, Phys. Rev. **D59**, 075003 (1999).

[6] J. F. Gunion, H. E. Haber, G. Kane and S. Dawson, *"The Higgs Hunter's Guide"*, Addison-Wesley, 1996.

[7] See, *e.g.* C. Caso *et al.* [PDG], Eur. Phys. J. **C3**, 1 (1998).

[8] H. Cheng, B. A. Dobrescu and K. T. Matchev, Phys. Lett. **B439**, 301 (1998); Nucl. Phys. **B543**, 47 (1999).

[9] H. Cheng, B. A. Dobrescu and C. T. Hill, hep-ph/9912343.

[10] G. Rupak and E. H. Simmons, Phys. Lett. **B362**, 155 (1995); S. Raychaudhuri and A. Raychaudhuri, Phys. Rev. **D44** (1991) 2663.

[11] A. Datta, A. Raychaudhuri, S. Raychaudhuri and S. Chakrabarti, Z. Phys. **C72**, 449 (1996).

[12] J. Blumlein *et al.*, Z. Phys. **C51**, 341 (1991).

[13] A. Datta and A. Raychaudhuri, Phys. Rev. **D57**, 2940 (1998).

[14] P. Abreu *et al.* [DELPHI Collaboration], Phys. Lett. **B458**, 431 (1999); G. Abbiendi *et al.* [OPAL Collaboration], Phys. Lett. **B464**, 311 (1999).

[15] B. Abbott *et al.* [D0 Collaboration], Phys. Rev. Lett. **82**, 2244 (1999); P. J. Wilson [CDF Collaboration], preprint FERMILAB-Conf-98-213-E.

[16] See, *e.g.* S. P. Martin and J. D. Wells, Phys. Rev. **D60**, 035006 (1999) and references therein.

[17] G. Landsberg and K. T. Matchev, hep-ex/0001007.

[18] T. Sjostrand, Comput. Phys. Commun. **82**, 74 (1994).

[19] J. Conway, talk given at the SUSY/Higgs Workshop, Fermilab, May 14-16, 1998; additional information at www.physics.rutgers.edu/~jconway/soft/shw/shw.html.

[20] J. D. Lykken and K. T. Matchev, Phys. Rev. **D61**, 015001 (2000); hep-ex/9910033; K. T. Matchev and D. M. Pierce, Phys. Rev. **D60**, 075004 (1999).

[21] A. Pukhov *et al.*, hep-ph/9908288.

[22] S. Mrenna and J. Wells, hep-ph/0001226.