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Probing the nature of the Higgs-like boson via $h \rightarrow V F$ decays

Gino Isidori $^{a,b}$, Aneesh V. Manohar $^c$, Michael Trott $^{b,*}$

$^a$ INFN, Laboratori Nazionali di Frascati, I-00044 Frascati, Italy
$^b$ Theory Division, Physics Department, CERN, CH-1211 Geneva 23, Switzerland
$^c$ Department of Physics, University of California at San Diego, La Jolla, CA 92093, United States

**A R T I C L E I N F O**

**A B S T R A C T**

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1. Introduction

Characterizing the properties of the newly discovered scalar boson at the LHC [1] is of central importance to determine experimentally the nature of electroweak symmetry breaking, and to investigate the possibility of physics beyond the Standard Model (SM). It is particularly important to determine how the new boson couples to the $SU(2)_L \times U(1)_Y$ gauge fields, since these couplings are directly related to symmetry breaking and gauge boson mass generation. The 125 GeV boson cannot decay into an on-shell pair of massive gauge bosons, but it can decay via $h \rightarrow V V^*$, $V^* \rightarrow F$, where one of the bosons is off-shell.

The purpose of this Letter is to show that the offshellness of $V^*$ can be viewed as a virtue in Higgs studies. It allows one to measure decays to final states $F$ that would not be accessible if on-shell decays were allowed, and the additional decay channels increase the sensitivity to new-physics (NP) effects, as they affect kinematic distributions of $F$ in addition to the total rate. We demonstrate this conclusion using two examples: (a) $F$ is a pair of light leptons $\ell^{+}\ell^{-}$ or $\ell \nu$, with $\ell = e, \mu$, and (b) $F = P$ is a hadronic state composed of a single pseudoscalar or vector meson.

In the two-lepton case, most of the interesting information is encoded in the two-dimensional kinematic distributions of the leptons in the Higgs rest frame. We analyze such distributions both in the SM, and in the context of a general effective field theory (EFT) approach, neglecting lepton masses. We show that these distributions, which will soon be accessible at the LHC with increasing statistics, contain information that cannot be directly extracted from a global fit to Higgs signal strengths. In the $h \rightarrow VP$ case we show how these rare processes, with SM rates in the $10^{-5}$ range, can provide a complementary tool to extract Higgs properties not easily accessible from the purely leptonic modes.

2. Amplitude decomposition

Consider the $h \rightarrow V F$ amplitude where $V = [W^{\pm}, Z^0]$ is an on-shell massive weak gauge boson while $F$ is a final state generated at tree level by the electroweak charged or neutral currents,

$$\mathcal{L}_I = \frac{e}{\sqrt{2} \sin \theta_W} J_\mu^V W_{\pm}^\mu + \frac{e}{\sin \theta_W \cos \theta_W} J_\mu^F Z^\mu. \quad (1)$$

Let $f_{\mu}^{F,V} = (\mathcal{F} | J_{\mu}^{V,F}) | 0 \rangle$ be the matrix element relevant for $V \rightarrow F$. The decay amplitude $A[h \rightarrow V (\ell, p) F(q)]$ can be decomposed in terms of four independent Lorentz structures, which we define as

$$A_F = C_V g_2^2 m_\nu \left\{ \begin{array}{c} \frac{g_\ell g_2}{q^2 - m_\nu^2} \left[ f_1^V (q^2) g_{\mu\nu} + f_2^V (q^2) q_\mu q_\nu \right. \\
+ f_3^V (q^2) (p \cdot q g_{\mu\nu} - q_\mu p_\nu) \left. \right]
+ f_4^V (q^2) \epsilon_{\mu\nu\rho\sigma} p_\rho q_\sigma. \quad (2) \end{array} \right.$$
used \( p \cdot \epsilon = 0 \) for physical \( V \) bosons but we have not made any assumption about the angular momentum of the \( \mathcal{F} \) state. We will also use the dimensionless variables \( \rho = m_\ell^2/m_\ell^2, \hat{q}^\mu = q^\mu/m_\ell, \hat{p}^\mu = p^\mu/m_\ell \). \( f_1^{\mathcal{F}} \) are real, and \( f_2^{\mathcal{F} \mu \nu} \) are complex conjugates of each other. For a \( 0^+ \) scalar \( h \), \( f_1^{W \mu \nu} \) and \( f_2^{W \mu \nu} \) \( \neq 0 \) violate CP.

In general, the \( f_i(q^2) \) are four independent dimensionless form factors. At \( q^2 = m_\ell^2 \), two of them satisfy the relation

\[
f_1^{\mathcal{F}}(m_\ell^2) = -m_\ell^2 f_2^{\mathcal{F}}(m_\ell^2),
\]

dictated by the requirement that the pole of the amplitude when \( q^2 \to m_\ell^2 \) corresponds to the exchange of an off-shell \( V \). In the SM, \( f_1^{\mathcal{F} \mu \nu} = 1, f_2^{\mathcal{F} \mu \nu} = -1/m_\ell^2 \), and \( f_4^{\mathcal{F} \mu \nu} = 0 \).

The different form-factors can be probed by using different final states. The differential decay rate summing over \( V \) polarizations is

\[
d\Gamma = \frac{\pi^2 C_F^2 v^4 m_\ell^2}{2m_\ell} \cdot \frac{\mathcal{M}^{\mu \nu} f_1^{\mathcal{F}} f_2^{\mathcal{F} \mu \nu}}{(q^2 - m_\ell^2)^2} \cdot \lambda(q^2, \rho) \, dq^2 \, d\Phi_\mathcal{F},
\]

where \( \lambda(q^2, \rho) = \sqrt{(1 + q^2 - \rho^2 - 4q^2)} \).

The simplified form factors are

\[
f_1^{\mathcal{F}}(q^2) = c_1 + g_2(c_2 + c_3) \left( 1 + \frac{q^2}{m_\ell^2} \right),
\]

\[
f_2^{\mathcal{F}}(q^2) = -\frac{1}{m_\ell^2} \left[ c_1 + 2g_2(c_2 + c_3) \right],
\]

\[
f_4^{\mathcal{F}}(q^2) = \frac{2g_2^2}{m_\ell^2} c_3, \quad f_3^{\mathcal{F}}(q^2) = 0.
\]

The three parity-conserving form factors correspond to three independent combinations of the parameters of the EFT Lagrangian, and only one combination is determined by the total decay rate. The dependence of the differential rate on \( f_i(q^2) \) in Eqs. (6)–(7) offers the opportunity to determine the individual form-factors, and hence the independent operator coefficients with sufficient data.

4. \( \mathcal{F} = \ell^+ \ell^- \), \( \ell \nu \)

There are two kinematic variables needed to describe the final state, after averaging over lepton spins. Two convenient choices are either the two lepton energies in the \( h \) rest frame \((E_{1,2})\), or the lepton invariant mass \( q^2 \) and \( c_0 = \cos \theta \), where \( \theta \) is the angle between the lepton axis and the dilepton rest frame and the Higgs momentum. For these two cases we can write \((y_1 = 2E_i/m_\ell, \text{ with } i = 1 \text{ for the lepton and } i = 2 \text{ for the antilepton})\)

\[
\frac{d^2\Gamma}{dy_1 \, dy_2} = \frac{2m_\ell^2}{\lambda(\hat{q}^2, \rho)} \frac{d^2\Gamma}{dq^2 \, dc_0} = \frac{C_\ell^2 m_\ell^4 m_\ell^2 m_\ell}{256\pi^4} \left[ \mathcal{M}_i^{\mu \nu} f_1^{\mathcal{F}} f_2^{\mathcal{F} \mu \nu} \right] \left( \frac{q^2}{m_\ell^2} \right)^2.
\]

Neglecting lepton masses, the term between square brackets can be evaluated from Eq. (6) using

\[
\chi_c \sum_{\ell \text{ spins}} J \cdot J = -2\hat{q}^2 = 2(1\rho - y_1 - y_2),
\]

\[
\chi_c \sum_{\ell \text{ spins}} (p \cdot J)(p \cdot J^\dagger) = -\hat{q}^2 + y_1 y_2 = \frac{1}{4} \lambda^2(\hat{q}^2, \rho)(1 - c_0^2),
\]

where \( \chi_c = (g_0^2)^2 + (g_1^2)^2 \), \( Y_c = (g_0^2)^2 - (g_1^2)^2 \).

The general expression of the double differential energy distribution can be deduced from Eqs. (4)–(11).
\[
\frac{d^2 \Gamma}{dq^2 \, dc_\ell} = \frac{C_q^2 q^2 \lambda(q^2, \rho)}{256 \pi^3 m_h \, (q^2 - m_h^2)^2} \times \left\{ X_c q^2 \left[ f_1 + \frac{1}{2} (m_h^2 - 2 \rho - m_h^2) f_3 \right]^2 + \frac{1}{4} m_h^2 \lambda(q^2, \rho) f_4 a_1 \right\} \\
+ \frac{1}{8 X_c m_h^2 \lambda(q^2, \rho)} \left[ \frac{f_1 v_c^2}{m_V^2} - q^2 f_3 s^2 - q^2 f_4 a_1 \right] (1 - c_\ell^2) \\
- Y_c \text{Im} \left( f_1^* + \frac{1}{2} (m_h^2 - 2 \rho - m_h^2) f_3 \right) f_4 m_h^2 q^2 \lambda(q^2, \rho) c_\ell \right\},
\]

(12)

with \(0 \leq q^2 \leq (1 - \sqrt{s})^2\) and \(-1 \leq c_\ell \leq 1\). The \(q^2\) spectrum is particularly simple and sensitive to possible deviations from SM.

In the SM, \(f_1 = 1, f_3 = 0\), and the double differential rate is

\[
\frac{1}{N_{SM}} \frac{d^2 \Gamma^{SM}}{dy_1 \, dy_2} = \frac{\rho^2 + 2 (y_1 + y_2 - \frac{3}{2}) + \frac{1}{2} (1 - y_1)(1 - y_2)}{y_1 - y_2^2}
\]

(13)

where \(N_{SM} = X_c C_q^2 q^2 m_h/(128 \pi^4)\) and \(0 \leq y_1 \leq (1 - \rho)\) and \((1 - \rho - y_1) \leq y_2 \leq (1 - \rho - y_1)/(1 - y_1)\). The lepton energy spectrum is

\[
\frac{2}{N_{SM}} \frac{d \Gamma^{SM}}{dy_1} = \frac{y_1(1 - y_1 - \rho)}{\rho(1 - y_1)} \left[ 2\rho^2 - \rho + y_1(1 - y_1) \right] \log \left[ \frac{\rho - y_1(1 - y_1)}{(1 - y_1)} \right].
\]

(14)

The usefulness of these differential distributions is illustrated in Fig. 1, where we compare different spectra, with \(c_i\) chosen to leave the total rate unchanged. \(1\) The \(d^2 \Gamma/dq^2\) spectrum exhibits large shape variations, which can be directly probed experimentally, and used to constrain the EFT. The variation in the \(q^2\) spectrum is due to a modified weighting of the terms in Eq. (11), which have a different \(q^2\) dependence.

On the other hand, the dependence of the lepton energy spectrum \(d \Gamma/dy_1\) on \(c_i\) is much weaker. Integrating over \(y_2\) averages over a wide range of \(q^2\). As a result the shape of the \(d \Gamma/dy_1\) distribution is almost universal, even in presence of the NP effects in the EFT we consider. This stability of \(d \Gamma/dy_1\) does not make it uninteresting — it provides a check for reconstruction errors or unaccounted for experimental systematics. The area under the curve of this spectrum is the total decay rate, and deviates from the SM value while maintaining this common shape.

We have examined the possibility of using moments of the lepton energy and \(q^2\) spectra to extract the Wilson coefficients of the operators in the EFT. However, these moments depend weakly on the \(c_i\). A more promising observable is the asymmetry \(A^{q^2}\) (either differential or integrated) given by weighting \(d^2 \Gamma/dy_1 \, dy_2\) by \(\chi = (y_2 - y_1)/(1 - y_1/2)\) (1 - \(y_1\)) is the midpoint of the \(y_2\) integration range. \(A^{q^2}\) is very sensitive to modifications of the form factors. In the SM, the integrated asymmetry is \(A^{q^2} = 15\%\), but it ranges from \(-9\%\) to \(+27\%\) for the illustrative EFT parameters adopted in Fig. 1.

5. Mesonic decays

The \(h \to VP\) amplitude depends on the current matrix element

\[
\langle P(q)|J_\mu(0)|0\rangle = \frac{1}{2} F_P q^\mu,
\]

(15)

where \(F_P\) is the pseudoscalar meson decay constant. This current selects a single form-factor combination

\[
f_F^V(q^2) = \frac{f_F^V(q^2) + q^2 f_F^V(q^2)}{q^2 - m_V^2},
\]

(16)

which has no pole as \(q^2 \to m_V^2\), by Eq. (3). In the SM, one has

\[
g_V^2 f_P^{VP,SM}(q^2) = -\frac{g_V^2}{m_V^2} = -\frac{4}{v^2},
\]

(17)

which is independent of \(g_V\) (here \(v = 246\) GeV). It is instructive to look at the structure of the amplitude for the \(h \to VP\) process,

\[
(A_P^{SM})^2 = \frac{-g_V^2 C_V F_P}{4 v} (\varepsilon \cdot q).
\]

(18)

This amplitude probes the ratio of two order parameters which both break \(SU(2) \times U(1)\) in the SM, \(F_P\) from the quark condensate and \(v\) from the Higgs sector. It provides a very interesting probe of the vacuum structure of the theory.\(^2\)

Compared to the leading decay mode of a light Higgs, \(h \to bb\), the \(h \to VP\) decay amplitude is parametrically suppressed by the small ratio \(F_P/m_b\). In the limit where we neglect the mass of the pseudoscalar meson,

\[
\Gamma(h \to VP)^{SM} = \frac{m_h^2}{6 v^2} |C_V F_P| \left( \frac{m_V}{m_h} \right)^3
\]

(19)

where \(\Gamma(h \to bb)^{SM} = 3m_h m_b^2/(8\pi v^2)\). This expression holds both for \(V = Z\) and \(V = W^\pm\), separately for each sign of charge.

Given the normalization of the currents in Eq. (1), the explicit expressions of \(F_P\) in some of the most interesting modes are \(F_K = V_{ud} f_\pi\), \(F_{KK} = V_{us} f_\pi\), \(F_{\pi 0} = f_\pi/\sqrt{2}\), \(F_{D0} = V_{ud} f_0\), \(F_{D^0} = V_{us} f_{D^0}\), \(F_{b_0} = V_{cb} f_{D^0}\), and \(F_{b_0} = f_{b_0}/2\), where \(F_P\) are the standard meson decay constants \(f_\pi \approx 130\) MeV, \(f_K \approx 160\) MeV, \(f_0 \approx 207\) MeV, \(f_{D^0} \approx 250\) MeV, and \(f_{b_0} \approx 400\) MeV [13–15]. From these values we deduce the SM rates reported in Table 1. Despite the smallness of these rates, and the huge background at the LHC, we stress that some channels may have a relatively clean experimental signature, due to the displaced vertex of the subsequent \(P\) decay.

Within the general EFT approach the \(h \to VP\) rate assumes the following form

\[
\frac{\Gamma(h \to VP)}{\Gamma(h \to VP)^{SM}} = \frac{m_h^2}{6 v^2} |C_V F_P| \left( \frac{m_V}{m_h} \right)^3 \left[ c_1 + g_V^2 (c_2 + c_3) \right]^2,
\]

(20)

with possible \(O(1)\) variations with respect to the SM. These variations are closely related to the possible variation of the \(d^2 \Gamma(h \to VP)/dq^2\) spectrum at \(q^2 = 0\), which is quite difficult to access experimentally. As a further illustration of the complementarity of \(h \to VV\) and \(h \to VP\) modes, we report here the dependence of the two total rates from the EFT parameters, adopting a common normalization for the leading term:

\(^2\) The combination \(F_P/v\) appears also in the purely leptonic \(P \to \ell\ell\) decays; however, in that case it is a probe of the Goldstone component of the Yukawa interaction (as manifest from the gaugeless limit of the theory, see e.g. [12]). Computing the \(h \to VP\) amplitude in the \(g \to 0\) limit treating \(V\) as an external field shows that the \(h \to VP\) amplitude probes the trilinear \(h \gamma_{\mu} \psi V^\mu\) coupling, where \(\psi\) is an (eaten) Goldstone boson.
\begin{equation}
\Gamma_{\ell\ell} \propto c_1^2 + 0.9 c_1 c_2 + 1.3 c_1 c_3 + 0.6 c_2 c_3 + 0.2 c_3^2 + 0.5 c_3^2, \\
\Gamma_{VP} \propto c_1^2 + 0.9 c_1 c_2 + 0.9 c_1 c_3 + 0.4 c_2 c_3 + 0.2 c_2^2 + 0.2 c_3^2.
\end{equation}

In the limit where we neglect light hadron masses, Eqs. (19) and (20) continue to hold with \( P \to P^* \), where \( P^* \) is a vector meson, with decay constant defined by
\begin{equation}
\langle P^*(q, \epsilon) | J_\mu | 0 \rangle = \frac{1}{2} F_P m_P \epsilon^\mu.
\end{equation}

The corresponding SM rates are given in Table 1 using \( f_0 \approx 157 \text{ MeV} \), \( f_\rho \approx 210 \text{ MeV} \), \( f_{J/\psi} \approx 410 \text{ MeV} \) [13,17] and \( f_{D_{10}^-} / f_{D_{10}^+} \approx 1.3 \) [16]. The \( P^* \) is longitudinally polarized. For heavy quarks, spin symmetry implies \( f_P = f_{P^*} \) [18]; the vector structure of \( f_{J/\psi} \) implies \( F_{\rho^0} = (1 - 2 s_W^2) f_\rho / \sqrt{2} \), \( F_\phi = (1/2 - 2/3 s_W^2) f_\phi \), \( F_{J/\psi} = (1 - 4/3 s_W^2) f_{J/\psi} \).

6. Conclusions

We have shown the importance and utility of a decomposition of the \( h \to V F \) amplitude into form-factors which can be probed by different final states, and how differential spectra can be used to disentangle the \( h V V^* \) couplings in a general EFT approach. Complementary information is provided by leptonic and exclusive (semi)-hadronic final states, among which the \( h \to VP \) decay is a simple and particularly interesting example. See [19] for a related analysis in the associated production process.
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