Anomalous Higgs Boson Production and Decay

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ABSTRACT

A non-standard symmetry breaking sector may lead to derivative couplings of the Higgs boson and thereby to anomalous interactions with the electroweak gauge bosons. In the framework of gauge-invariant effective Lagrangians the resulting changes in Higgs boson production and decay mechanisms are related to anomalous triple vector boson couplings. Using low energy constraints on the dimension-six operators in the effective Lagrangian, we discuss the size of deviations from SM predictions which may be expected in Higgs boson production and decay rates. Large enhancements are allowed \textit{e.g.} in the $Z \to H\gamma$ and $H \to \gamma\gamma$ partial decay widths, leading to $Z \to \gamma\gamma\gamma$ events at LEP.
In recent years the standard model (SM) of electroweak interactions has been beautifully confirmed, in particular via the LEP experiments. $Z$ production via $e^+e^-$ annihilation and the bulk of the $Z$ decay processes mainly probe the couplings of the quarks and leptons to the $Z$ boson, however. While the gauge theory predictions of these fermion-$Z$ couplings have now been confirmed at the 1% level or better, the bosonic sector of the SM has been tested to a much lesser degree because present accelerators largely operate below weak boson pair production or Higgs production threshold.

In order to find out how the $SU(2) \otimes U(1)$ symmetry of the SM is broken in nature, an experimental determination is needed of the interactions between the gauge bosons and the remnants of the order parameter which gives rise to the spontaneous breaking of the gauge symmetry, i.e. the Higgs boson in the SM. In this letter we investigate the phenomenology of models which are relatively close to the SM in that they possess a (possibly light) Higgs scalar as the remnant of the $SU(2)$ doublet order parameter.

The anomalous interactions of this doublet field $\Phi$, which possesses the same quantum numbers as the SM Higgs doublet field, can conveniently be described by an effective Lagrangian

\[ L_{\text{eff}} = \sum_i \frac{f_i}{\Lambda^2} O_i + \sum_i \frac{f_i^{(8)}}{\Lambda^4} O_i^{(8)} + \ldots. \]

(0.1)

Here the scale $\Lambda$ may be identified with the typical mass of new particles associated with the fundamental interactions underlying the symmetry breaking sector. We here assume that the $W$ and the $Z$ are indeed gauge bosons of an $SU(2) \otimes U(1)$ local symmetry. Derivative couplings of the Higgs, as described by some of the operators in $L_{\text{eff}}$, are then related to Higgs-gauge boson (e.g. $HVV$) interactions which lead to new phenomena like enhanced $H \rightarrow \gamma\gamma$ decay rates or $Z \rightarrow H\gamma$ production. The purpose of this letter is to relate the expected/possible size of such effects to bounds [1] derived from present low energy data and to the measurement of anomalous triple gauge boson vertices (TGV’s).

For our analysis it is sufficient to consider the dimension six operators in $L_{\text{eff}}$ only. This allows a qualitative analysis which is quite model independent and which is quantitatively correct if $m_H << \Lambda$ and $v << \Lambda$, where $v$ is the vacuum expectation value of the Higgs doublet field. A complete analysis of all dimension six operators has been presented by Buchmüller and Wyler [2]. Here it suffices to consider operators which can be constructed...
out of the Higgs field $\Phi$, covariant derivatives of the Higgs field, $D_\mu \Phi$, and the field strength tensors $W_{\mu\nu}$ and $B_{\mu\nu}$ of the $W$ and the $B$ gauge fields:

$$[D_\mu, D_\nu] = \hat{B}_{\mu\nu} + \hat{W}_{\mu\nu} = i \frac{g'}{2} B_{\mu\nu} + i g \frac{\sigma^a}{2} W^a_{\mu\nu}. \quad (0.2)$$

For our discussion of non-standard HVV couplings six such operators need to be considered, which we call $O_{\Phi,1}$, $O_{BW}$, $O_W$, $O_B$, $O_{WW}$, and $O_{BB}$. In addition the operator $O_{\Phi,2}$ contributes via the Higgs boson wave function renormalization. They are given explicitly by

$$L_{\text{eff}} = \frac{7}{\Lambda^2} \sum_{i=1}^7 \frac{f_i}{\Lambda^2} O_i = \frac{1}{\Lambda^2} \left( f_{\Phi,1} (D_\mu \Phi)^\dagger \Phi \Phi^\dagger (D^\mu \Phi) 
+ \frac{1}{2} f_{\Phi,2} \partial_\mu (\Phi^\dagger \Phi) \partial^\mu (\Phi^\dagger \Phi) + f_{BW} \Phi^\dagger \hat{B}_{\mu\nu} \hat{W}^{\mu\nu} \Phi 
+ f_W (D_\mu \Phi)^\dagger \hat{W}^{\mu\nu} (D_\nu \Phi) + f_B (D_\mu \Phi)^\dagger \hat{B}^{\mu\nu} (D_\nu \Phi) 
+ f_{WW} \Phi^\dagger \hat{W}_{\mu\nu} \hat{W}^{\mu\nu} \Phi + f_{BB} \Phi^\dagger \hat{B}_{\mu\nu} \hat{B}^{\mu\nu} \Phi \right). \quad (0.3)$$

When the Higgs doublet field is replaced by its v.e.v. one finds that the operators $O_{\Phi,1}$, and $O_{BW}$ contribute to gauge boson self-energies. $O_{\Phi,1}$ changes the $Z$ but not the $W$ mass at the tree level and hence is severely constrained by the measured small value of $\delta \rho$. $O_{BW}$ leads to $W^3$–$B$ mixing and hence contributes to the $S$ parameter of Peskin and Takeuchi [3]. A recent analysis of low energy constraints found

$$f_{\Phi,1} \Lambda^2 = (0.11 \pm 0.20) \text{ TeV}^{-2}, \quad f_{BW} \Lambda^2 = (1.9 \pm 2.9) \text{ TeV}^{-2}. \quad (0.4)$$

Similarly the operators $O_W$ and $O_B$ give rise to anomalous TGV’s

$$\kappa_\gamma = 1 + (f_B + f_W) \frac{m^2_W}{2\Lambda^2}, \quad \kappa_Z = 1 + (c^2 f_W - s^2 f_B) \frac{m^2_Z}{2\Lambda^2}, \quad g_1^Z = 1 + f_W \frac{m^2_Z}{2\Lambda^2}, \quad g_1^W = 1 + \frac{m^2_W}{2\Lambda^2}, \quad (0.5)$$

in the notation of the phenomenological Lagrangian

$$i L_{\text{eff}}^{WW} = g_{WWV} \left( g_1^V (W_{\mu\nu} W^\mu \mu - W^\dagger W_{\mu\nu}) V^\nu + \kappa_V W^\dagger W_{\mu\nu} V^\nu \right), \quad (0.6)$$

with the overall coupling constants defined as $g_{WW\gamma} = e$ and $g_{WWZ} = e \cot \theta_W$. Within the SM the couplings are given by $g_1^Z = g_1^V = \kappa_Z = \kappa_\gamma = 1$. A direct measurement of the $WW\gamma$ vertex exists from $W\gamma$ production at hadron colliders with the result $\kappa_\gamma = 1^{+2.6}_{-2.2}$ which translates into a measurement of $f_B + f_W$.
A more stringent bound can be obtained from an analysis of 1-loop effects of anomalous TGV’s on low energy observables. For $m_H = 200$ GeV and a top quark mass of 140 GeV, one obtains \[ f_B + f_W = 0^{+800}_{-700} \text{ TeV}^{-2}. \] (0.7)

These bounds neglect possible correlations or cancellations between various new physics contributions, however, and can only be considered as order of magnitude estimates. The same is true for the coefficients of the operators $O_{WW}$ and $O_{BB}$, which can only be constrained via their 1-loop contributions to low energy observables prior to Higgs discovery. For the same assumptions as made for Eq. (0.8) one finds \[ \frac{f_W}{\Lambda^2} = (3 \pm 27) \text{ TeV}^{-2}, \quad \frac{f_B}{\Lambda^2} = (13 \pm 19) \text{ TeV}^{-2}. \] (0.8)

(0.9)

Allowing for modest correlations/cancellations among the various operators in their effects on low energy observables, the coefficients of the four operators $O_W, O_B, O_{WW}$, and $O_{BB}$ may be as large as 100 TeV$^{-2}$, while for the operators $O_{\Phi,1}$ and $O_{BW}$ an upper bound compatible with the low energy data is of order $|f_i|/\Lambda^2 = 1$ TeV$^{-2}$. We shall take these values as illustrative examples in the following to estimate the size of effects in Higgs boson production and decay. Notice that values $|f_B|/\Lambda^2, |f_W|/\Lambda^2 = 100$ TeV$^{-2}$ correspond to anomalous TGV’s of order $\kappa_\gamma - 1 = 0.3 \ldots 0.6$ and hence are exactly in the interesting range for $W^+W^-$ production experiments at LEP II \[ \text{[7]} \]. De Rújula and collaborators \[ \text{[4]} \] have argued that it is unnatural to expect a factor of 100 difference between the coefficients of the various operators. We may rather take the more stringent constraints of Eq. (0.4) as an estimate of the bounds for all the operators in the effective Lagrangian. Following this more conservative view we shall consider the case when all coefficients are of order $|f_i|/\Lambda^2 = 1$ TeV$^{-2}$ as a second illustrative example for the consequences of the anomalous HVV couplings.

The $H\gamma\gamma$, $HZZ$, $H\gamma Z$, and $HWW$ couplings follow from the effective Lagrangian (0.1) by making the replacement $\Phi \rightarrow (0, (v + H)/\sqrt{2})^T$. For later use we here list the results in terms of the HVV effective Lagrangian,
The total width for this process is given by

$$
\Gamma(H \to \gamma\gamma) = \frac{\alpha s^2 m_W^2 m_H^3}{4} \left| \frac{f_{BB} + f_{WW} - f_{BW}}{\Lambda^2} \right|^2 + \frac{\alpha}{8\pi s^2 m_W^2} I^2 ,
$$

where $g^2 = e^2/s^2 = 8m_W^2 G_F/\sqrt{2}$.

In addition the operators $O_{\Phi,1}$ and $O_{\Phi,2}$ renormalize the weak boson masses and the Higgs boson wave function. By using the observed masses and couplings the HVV interactions are modified as

$$
\mathcal{L}_{\text{eff}}^{\text{HVV}} = g m_W \left\{ -\frac{s^2 (f_{BB} + f_{WW} - f_{BW})}{2} H A_{\mu\nu} A^{\mu\nu} + \frac{2 m_W^2 f_{\Phi,1}}{g^2} H Z_{\mu\nu} Z^\mu Z^\nu \\
+ \frac{c^2 f_W + s^2 f_B Z_{\mu\nu} (\partial^\nu H) - s^2 f_{BB} + c^2 f_{WW} + s^2 c^2 f_{BW} H Z_{\mu\nu} Z^\mu}{2c^2} \\
+ \frac{s (f_W - f_B)}{2c} A_{\mu\nu} Z^\mu (\partial^\nu H) + \frac{s (2s^2 f_{BB} - 2c^2 f_{WW} + (c^2 - s^2) f_{BW})}{2c} H A_{\mu\nu} Z^\mu \\
+ \frac{f_W}{2} (W^+ W^- + W^- W^+) (\partial^\nu H) - f_{WW} W^\nu W^{-\nu} \right\} ,
$$

(0.10)

Here the couplings are normalized as $g^2 = g_{\gamma W}(0)$ and $g_{Z}^2 = g_{Z}^2(m_Z^2)$ of ref. [1], whose magnitudes have been precisely measured. Notice that the correction terms in $\mathcal{L}_{\text{eff}}^{\text{HVV}}$ need to be considered in addition to the terms given in Eq. (10). They are valid only for $|f_{\Phi,1}| v^2/\Lambda^2, |f_{\Phi,2}| v^2/\Lambda^2 \ll 1$. This condition clearly is satisfied for $f_{\Phi,1}$ (see Eq. (4)). The only effect of the operator $O_{\Phi,2}$ is a finite wave function renormalization of the Higgs field by a factor $Z_{H}^{W} = 1/\sqrt{1 + f_{\Phi,2} v^2/\Lambda^2}$. The phenomenological consequence is a common rescaling of all Higgs production rates and partial decay widths by a factor $Z_{H}$. We shall mostly neglect this overall factor (by setting $f_{\Phi,2} = 0$) when considering the nonstandard contributions to $H \to VV$ decay rates and $Z \to HV$ decay, which are induced by the anomalous HVV couplings.

**Higgs Decays.** Higgs decays into $\gamma\gamma$, $Z\gamma$, $ZZ$, and $W^+W^-$ are affected by non-standard interactions from dimension-six operators. In the SM the decay $H \to \gamma\gamma$ occurs at the one-loop level. All massive particles with non-zero electromagnetic charge run through the loop, the most important contributions arising from the top quark and the $W$ boson. However, dimension-six operators contribute at the tree level and can therefore lead to large deviations from SM expectations [8]. The total width for this process is given by

$$
\Gamma(H \to \gamma\gamma) = \frac{\alpha s^2 m_W^2 m_H^3}{4} \left| \frac{f_{BB} + f_{WW} - f_{BW}}{\Lambda^2} \right|^2 + \frac{\alpha}{8\pi s^2 m_W^2} I^2 ,
$$

(0.12)
with $s = \sin \theta_W$. The SM contribution is parametrized by the complex-valued function 

$$I = \sum_i N_{ci} e_i^2 F_i$$

where $N_{ci}$ is the color multiplicity of particle $i$ and $e_i$ is its charge. The functions $F_i$ are given explicitly in Ref. [9].

In the SM the branching fraction for $H \to \gamma\gamma$ is $\mathcal{O}(10^{-3})$. Although it is a rare decay it is the primary search mode for an intermediate-mass Higgs boson at hadron colliders where the $H \to b\bar{b}$ signal is swamped by QCD backgrounds, but $H \to WW, ZZ$ is not kinematically allowed. For maximal values of $f_{BB}$ and $f_{WW}$ each ($\sim \mathcal{O}(100 \text{ TeV}^{-2})$), and assuming no large cancellations, the width for $H \to \gamma\gamma$ receives an enhancement by a factor $10^4$ and, below the $H \to WW^*$ “threshold”, is the primary decay mode, see fig. 1a. The enhanced width is shown by the dashed line, and the SM prediction is given by the solid line. For a heavier Higgs it may still be a competitive process with a branching fraction of a few percent. Hence, the search for an intermediate mass Higgs boson at hadron colliders would be greatly facilitated. Note, however, that a large $H\gamma\gamma$ coupling and large TGV’s occur in orthogonal directions in the $f_i$ parameter space. Hence a large rate for $H \to \gamma\gamma$ does not necessarily imply large TGV’s.

For “natural” values of the $f_i$ ($f_i/\Lambda^2 = 1\text{TeV}^{-2}$ for all $f_i$) the SM one-loop contribution and the dimension-six contributions are comparable; significant interference is, in principle, possible. Generically one might expect a change in the rate for $H \to \gamma\gamma$ by a factor of 2 or so. The dotted line displays destructive interference while the double-dotted line shows constructive interference.

The process $H \to \gamma Z$ is very similar to the process $H \to \gamma\gamma$; the SM contribution is a one-loop process, but the dimension-six contribution occurs at the tree level. The total width for this process is given by

$$\Gamma(H \to \gamma Z) = \frac{\alpha(m_H^2 - m_Z^2)^3 m_Z^2}{16 m_H^3} \left| f_W - f_B + 4s^2 f_{BB} - 4c^2 f_{WW} + 2(c^2 - s^2) f_{BW} \right| \frac{\alpha}{2 \pi s c m_Z^2} A^2.$$ (0.13)

The SM contribution is parametrized by the complex-valued function $A = A_F + A_W$ which is given explicitly in Ref. [9].

Setting $f_i/\Lambda^2 = 100\text{TeV}^{-2}$ for all $f_i$ except $f_{BW}$ then the width for $H \to \gamma Z$ is enhanced by a factor of $10^3$. See the dashed line in fig. 1b. The resulting branching fraction is around
10\% when the decay is kinematically allowed. Hence \( H \to \gamma Z \) would be a complementary channel to \( H \to \gamma \gamma \) for \( 100 \text{GeV} \leq m_H \leq 140 \text{GeV} \). For a heavier Higgs boson both \( H \to \gamma Z \) and \( H \to \gamma \gamma \) would complement \( H \to WW, ZZ \). The appearance of \( f_{BB} \) and \( f_{WW} \) in Eq. (1.1) means that one can have an enhanced \( H \to \gamma Z \) rate without large anomalous TGV’s. However, the appearance of \( f_B \) and \( f_W \) in Eq. (1.13) implies that, barring large accidental cancellations, measurable values for \( \Delta \kappa_\gamma, \Delta \kappa_Z \) and \( g^Z_1 \) imply a strongly enhanced \( H \to \gamma Z \) rate. E.g. a value of \( \kappa_\gamma \approx 0.2 \), at the limit of observability in \( W^+W^- \) production at LEP II [7], implies values of \( (f_B + f_W)/\Lambda^2 \) in the vicinity of the values used for the dashed line in fig. 1b.

For more “natural” values of the \( f_i \) (\( f_i/\Lambda^2 \sim \mathcal{O}(1 \text{TeV}^{-2}) \)) the SM one-loop contribution and the dimension-six contribution are comparable and interference is probable. One might expect a change in this rate by a factor of two or so, and this decay remains a rare process. See the dotted and double-dotted lines in fig. 1b. As such this channel is not likely to be instrumental in the discovery of the Higgs boson, but it does serve as an important precision test of the SM. Along with a measurement of the \( H \to \gamma \gamma \) rate and the measurement of anomalous TGV’s a measurement of the \( H \to \gamma Z \) rate places some important restrictions on the allowed directions in the \( f_i \) parameter space. As in the case of \( H \to \gamma \gamma \) there is a small region of parameter space corresponding to maximal destructive interference, in which case this mode would be unobservable.

The decays \( H \to ZZ \) and \( H \to WW \) occur at tree level, and hence they are affected significantly only when we allow larger magnitudes for the coefficients of dimension-six operators. Separating the result into longitudinal and transverse contributions the decay width for the process \( H \to ZZ \) is given by

\[
\Gamma(H \to ZZ) = \frac{\alpha}{128\pi} \frac{m_Z^3}{m_W^2} \sqrt{1 - x_z^2} \left[ x_z (1 + C_{ZZ}) + D_{ZZ} \right]^2,
\]

(0.14)

and

\[
\Gamma(H \to LL) = \frac{\alpha}{128\pi} \frac{m_L^3}{m_W^2} \sqrt{1 - x_z^2} \left[ (2 - x_z)(1 + C_{ZZ}) + D_{ZZ} \right]^2,
\]

(0.15)

where

\[
C_{ZZ} = (f_{\Phi,1} - f_{\Phi,2}) \frac{v^2}{2\Lambda^2} - 2 \frac{m_Z^2}{\Lambda^2} \left[ c^4 f_{WW} + s^2 c^2 f_{BW} + s^4 f_{BB} \right],
\]

\[
D_{ZZ} = 2 \frac{m_Z^2}{\Lambda^2} \left[ 2 c^4 f_{WW} + 2 s^2 c^2 f_{BW} + 2 s^4 f_{BB} - s^2 f_B - c^2 f_W \right],
\]
and $x_w = 4m_Z^2/m_H^2$.

The width for $H \to WW$ is very similar.

$$\Gamma(H \to W_TW_T) = \frac{g^2}{64\pi} \frac{m_H^3}{m_W^2} \sqrt{1 - x_w} \frac{1}{2} [x_w(1 + C_{WW}) + D_{WW}]^2,$$

(0.16)

and

$$\Gamma(H \to W_LW_L) = \frac{g^2}{64\pi} \frac{m_H^3}{m_W^2} \sqrt{1 - x_w} \frac{1}{4} [(2 - x_w)(1 + C_{WW}) + D_{WW}]^2,$$

(0.17)

where

$$C_{WW} = -(f_{\phi,1} + 2f_{\phi,2}) v^2 4\Lambda^2 - 2 m_W^2 f_{WW},$$

$$D_{WW} = 2 \frac{m_W^2}{\Lambda^2} (2f_{WW} - f_w),$$

and $x_w = 4m_Z^2/m_H^2$.

When anomalous contributions vanish, i.e. $f_i/\Lambda^2 \to 0$ for all $f_i$, then $C_{VV}, D_{VV} \to 0$ for $V = W, Z$ and the SM value for $\Gamma(H \to VV)$ is recovered. If $D_{VV} = 0$ then the $H \to VV$ decay width is enhanced by an overall factor $(1 + C_{VV})^2$. However, a non-zero value of $D_{VV}$ will change the ratio of longitudinally polarized and transversely polarized vector bosons. It is possible to construct a scenario where $C_{WW}, C_{ZZ} \neq 0$ with $D_{WW}, D_{ZZ} = 0$, but only in a very restricted region of parameter space. Hence, one should in general expect that a change in the overall rate will be accompanied by a change in the ratio of longitudinal and transverse polarizations.

The size of expected deviations in the $H \to WW$ and $H \to ZZ$ partial widths is shown in figs. 1c and 1d. For “natural” values of the $f_i$ ($f_i/\Lambda^2 \sim \mathcal{O}(1\text{TeV}^{-2})$) we do not expect measurable changes in these decay rates; the dotted and double-dotted lines are nearly degenerate with the solid SM curves. However, if the $f_i$ approach their phenomenological bounds then modest effects are likely to be seen. In this case destructive interference effects between SM and new physics contributions are possible, independently in the $H \to WW$ and $H \to ZZ$ decay modes. Hence the ratio of $WW$ vs. $ZZ$ Higgs signals at hadron supercolliders may be drastically altered.

**Z decays involving the Higgs boson.** The new interactions from dimension-six operators can significantly effect the $Z$ decay branching fractions $Z \to HZ^* \to Hf\bar{f}$, $Z \to H\gamma$
and \( Z \to H \gamma \to \gamma \gamma \gamma \). We present the differential decay rate for \( Z \to HZ^* \to Hf\bar{f} \) in terms of the individual particle momenta. Denote the four-momentum of the initial-state \( Z^0 \) by \( p_Z \). \( p_H \) denotes the four-momentum of the Higgs boson and \( E_H \) is its energy in the rest frame of the initial-state \( Z^0 \). Finally, \( q_1 \) and \( q_2 \) denote the fermion and anti-fermion four-momenta, and \( E_f \) is the energy of the fermion. Then

\[
\frac{d\Gamma}{dE_f dE_H} = \frac{g^4 m_W^2 g_1^2 + g_A^2}{24(2\pi)^3 e^6 m_Z^3 \left[ (p_Z - p_H)^2 - m_Z^2 \right]^2 + (m_Z \Gamma_Z)^2} \]

\[
\left\{ 2G_1^2 q_1 \cdot q_2 + (q_1 \cdot p_H q_2 \cdot p_H - m_H^2 q_1 \cdot q_2) \right. \\
\left. \frac{G_1^2}{m_Z^2} + 2G_1 G_2 \left( 1 - \frac{p_Z \cdot p_H}{m_Z^2} \right) - G_2^2 \left( m_H^2 - \frac{(p_Z \cdot p_H)^2}{m_Z^2} \right) \right\},
\]

(0.18)

where

\[
G_1 = -2 \left( 1 + (f_{\phi,1} - f_{\phi,2}) \frac{v^2}{2\Lambda^2} \right) + m_H^2 \frac{s^2 f_B + c^2 f_W}{\Lambda^2} \\
+ 4(m_Z^2 - p_Z \cdot p_H) \frac{s^4 f_{BB} + s^2 c^2 f_{BW} + c^4 f_{WW}}{\Lambda^2},
\]

(0.19)

\[
G_2 = 2 \frac{s^2 f_B + c^2 f_W}{\Lambda^2} - 4 \frac{s^4 f_{BB} + s^2 c^2 f_{BW} + c^4 f_{WW}}{\Lambda^2},
\]

(0.20)

and \( g_V = \frac{1}{2} T_3 - s^2 Q \) and \( g_A = -\frac{1}{2} T_3 \) are the vector and axial-vector couplings of the fermion.

For “natural” values of all \( f_i \) (\( f_i \sim \mathcal{O}(1\text{TeV}^{-2}) \)) new physics effects are negligible. For phenomenologically allowed values of \( f_i \sim \mathcal{O}(100\text{TeV}^{-2}) \) effects remain modest (see the dashed lines in fig. 2a.) but are important with regard to the ongoing Higgs boson search at LEP. Assuming SM couplings the Higgs boson is currently constrained to be heavier than approximately 60 GeV by searching in the channel \( Z \to HZ^* \to Hf\bar{f} \). Contributions from dimension-six interactions can weaken this bound significantly.

The decay \( Z \to H \gamma \) occurs in the SM at the one-loop level; the dimension-six contribution occurs at the tree level and has been considered previously by several authors [8]. The combined width for this process is

\[
\Gamma(Z \to H \gamma) = \frac{\alpha}{96 m_Z} \left( m_Z^2 - m_H^2 \right)^3 \\
\left| \frac{f_W - f_B + 4s^2 f_{BB} - 4c^2 f_{WW} + 2(c^2 - s^2) f_{BW}}{\Lambda^2} + \frac{\alpha}{2\pi s c m_Z^2} A \right|^2.
\]

(0.21)

The SM contribution is parametrized by the complex-valued function \( A \), which is given explicitly in Ref. [8]. If all \( f_i \) are of order \( f_i/\Lambda^2 \sim \mathcal{O}(1\text{TeV}^{-2}) \) then the SM and the dimension-six
contributions are comparable and, unless there is maximal destructive interference between the SM and the new contributions we do not expect a large change in the rate for \( Z \to H\gamma \). However, the contributions from dimension-six operators completely dominate this partial decay rate of the \( Z \) if the \( f_i \) are close to their phenomenological limits (see fig. 2b), and there is a large rate enhancement. The search for the Higgs boson in this channel actually provides a constraint on the linear combination of the \( f_i \) appearing in Eq. (??) (providing \( m_H < m_Z \)) since, as can be seen from the topmost curve (dashed line) in fig. 2b, for some allowed values of the \( f_i \) a light Higgs boson should have already been discovered. Notice that \( \Gamma(Z \to Hf\bar{f}) \) and \( \Gamma(Z \to H\gamma) \) involve different linear combinations of the \( f_i \), hence one process may be affected but not the other. In particular a reduced rate in the \( Z \to Hf\bar{f} \) channel does not necessarily imply an enhanced rate in the \( Z \to H\gamma \) rate, though an enhancement is likely.

Therefore, new physics of the type discussed here can weaken the lower limit on the mass of the Higgs boson obtained at LEP.

The sequential decay \( Z \to H\gamma \to \gamma\gamma\gamma \) is significant if both \( Z \to H\gamma \) and \( H \to \gamma\gamma \) branching fractions are enhanced by the new interactions. In the SM the decay \( Z \to \gamma\gamma\gamma \) occurs at the one loop level. The contribution due to fermions in the loop has been calculated \[10\] and, for a heavy top quark, the contribution is 0.7 eV. The contribution due to W bosons in the loop has also been calculated \[11\]. This contribution is found to be smaller by a factor of 35. Interference effects between bosonic and fermionic loops have not been calculated. We ignore the contribution from W bosons. \( \Gamma(Z \to \gamma\gamma\gamma) \) also receives a contribution via \( Z \to H\gamma \to \gamma\gamma\gamma \). In the purely SM scenario both the \( HZ\gamma \) vertex and the \( H\gamma\gamma \) vertex are generated at one loop, hence the SM contribution is non-negligible only if \( Z \to H\gamma \) is kinematically allowed.

With the inclusion of dimension-six effects both the \( HZ\gamma \) vertex and the \( H\gamma\gamma \) vertex may be enhanced. In this scenario the process \( Z \to H\gamma \to \gamma\gamma\gamma \) may be important even for a virtual Higgs boson. Because this process then involves the product of two dimension-six operators, our calculation via virtual Higgs boson exchange should be regarded as an estimate of the possible dimension-eight \( Z\gamma\gamma\gamma \) vertex. The results are summarised by fig. 3. The SM contribution is too small to be interesting. For small dimension-six contributions \( (f_i/\Lambda^2 \sim 1\text{TeV}^{-2}) \) the effects are also very small.
The situation is dramatically different for large dimension-six effects ($f_i/\Lambda^2 \sim 100\text{TeV}^{-2}$). Based upon fig. 3 the search for $Z \to \gamma\gamma\gamma$ events becomes promising for a light Higgs boson. There is even some hope for events in this channel for $m_H > m_Z$ if the Higgs boson is not too much heavier than $m_Z$.

**Summary.** New physics in the electroweak bosonic sector may be described by an effective Lagrangian of dimension-six operators. The coefficients of some of these operators are severely constrained by low-energy data ($f_i/\Lambda^2 < \mathcal{O}(1\text{TeV}^{-2})$), while others may be as large as $f_i/\Lambda^2 \sim \mathcal{O}(100\text{TeV}^{-2})$. Actually, allowing for arbitrary cancellations amongst the full set of operators no stringent and rigorous bounds exist on any of them.

In the pessimistic scenario ($f_i/\Lambda^2 < \mathcal{O}(1\text{TeV}^{-2})$ for all of the operators) one does not expect to see deviations from the SM predictions for TGV’s. Furthermore, one does not expect to observe changes in SM processes which occur at the tree level. However, processes which occur at one-loop in the SM but have tree-level dimension-six contributions could differ from their SM expectations appreciably. This is demonstrated in Fig. 4b where the various Higgs boson branching fractions are compared. While most of the branching fractions are indistinguishable from their SM values, the $H \to \gamma\gamma$ and $H \to Z\gamma$ rates are strongly affected, which would have important consequences for intermediate-mass Higgs boson searches at hadron colliders.

Huge effects on Higgs phenomenology are possible if large dimension six contributions close to their present phenomenological low energy bounds are realized in nature. This is demonstrated by the Higgs branching ratios of fig. 4a: an intermediate mass Higgs might predominantly decay into two photons. Actually large effects like the ones in fig.4a should be expected if anomalous TGV’s are large enough to be observed in $W^+W^-$ production at LEP II. Thus the search for a light Higgs at LEP I may have important consequences for vector boson pair production at higher energies.

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REFERENCES

[1] K. Hagiwara, S. Ishihara, R. Szalapski, and D. Zeppenfeld, preprint MAD/PH/737 (1993), Phys. Rev. D48, in press.

[2] W. Buchmüler and D. Wyler, Nucl. Phys. B268 (1986) 621.

[3] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. 65 (1990) 964.

[4] A. De Rújula, M. B. Gavela, P. Hernandez, E. Massó, Nucl. Phys. B384 (1992) 3.

[5] K. Hagiwara, K. Hikasa, R.D. Peccei, D. Zeppenfeld, Nucl. Phys. B282 (1987) 253.

[6] UA2 collaboration, J. Alitti et al., Phys. Lett. B277 (1992) 194; H. Aihara, talk given at the conference on Physics in Collisions, Heidelberg, June 1993.

[7] G. Barbiellini et al., in “Physics at LEP,” ed. by J. Ellis and R. D. Peccei, report CERN 86-02, Vol. 1; M. Davier et al., ECFA Workshop on LEP 200, ed. by A. Bohm and W. Hoogland, CERN report CERN 87-08, Vol. 1, p. 120.

[8] D. Düssedau and J. Wudka, Phys. Lett. 180B (1986) 290; L. Randall and N. Rius, Phys. Lett. B309 (1993) 365.

[9] J. F. Gunion, H. E. Haber, G.L. Kane, and S. Dawson, The Higgs Hunter’s Guide, Addison-Wesley, 1990.

[10] J. J. van der Bij and E. W. N. Glover, Nucl. Phys. B313 (1989) 237; See also J. J. van der Bij and E. W. N. Glover in, “Z Physics at LEP I”, ed. by G. Altarelli, R. Kleiss, and C. Verzegnassi, CERN report CERN 89-08, Vol. 2, p. 30; M.L. Laursen, K.O. Mikaelian and A. Samuel, Phys. Rev. D23 (1981) 2795.

[11] M. Baillargeon and F. Boudjema, Phys. Lett. B272 (1991) 158; Fang-xiao Dong, Xiang-dong Jiang, Xian-jian Zhou, Phys. Rev. D46 (1992) 5074.
FIGURES

FIG. 1. Higgs boson decay widths for the channels a) \( H \rightarrow \gamma\gamma \), b) \( H \rightarrow Z\gamma \), c) \( H \rightarrow WW \), d) \( H \rightarrow ZZ \). The solid line is purely SM, dots: \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for all six operators, double dots: \( f_i/\Lambda^2 = -1\text{TeV}^{-2} \) for all six operators, dashes: \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for \( f_{BW} \) and \( f_{\phi,1} \), while \( f_i/\Lambda^2 = 100\text{TeV}^{-2} \) for \( f_{BB}, f_{WW}, f_B, \) and \( f_W \), and long-dash short-dash: \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for \( f_{BW} \) and \( f_{\phi,1}, f_i/\Lambda^2 = -100\text{TeV}^{-2} \) for \( f_{BB}, f_{WW}, f_B, \) and \( f_W \).

FIG. 2. Partial \( Z \) decay widths with a Higgs boson in the final state. a) \( Z \rightarrow Hf\bar{f} \) summed over all kinematically allowed SM fermions, b) \( Z \rightarrow H\gamma \). The various lines are for the same linear combinations of dimension-six operators as in fig. 1: the solid line is purely SM, the dotted lines describe the effect of “phenomenologically allowed” coefficients \( f_i/\Lambda^2 \) while the dashed lines show examples of “natural” values of these coefficients.

FIG. 3. \( \Gamma(Z \rightarrow \gamma\gamma\gamma) \) in various scenarios. The solid line is purely SM, dots: \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for all six operators, double dots: \( f_i/\Lambda^2 = -1\text{TeV}^{-2} \) for all six operators and dashes: \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for \( f_{BW} \) and \( f_{\phi,1}, f_i/\Lambda^2 = 100\text{TeV}^{-2} \) for \( f_{BB}, f_{WW}, f_B \) and \( f_W \). The remaining curve (dash double dot) is the SM width for \( Z \rightarrow H\gamma \), included for reference.

FIG. 4. Higgs boson branching fractions for two choices of dimension-six operators: a) \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for \( f_{BW} \) and \( f_{\phi,1}, f_i/\Lambda^2 = 100\text{TeV}^{-2} \) for \( f_{BB}, f_{WW}, f_B \) and \( f_W \), b) \( f_i/\Lambda^2 = 1\text{TeV}^{-2} \) for all six operators.