Analyzing the *CP* Nature of a New Scalar Particle via $S \to Zh$ Decays

Martin Bauer,[1] Matthias Neubert,[2,3] and Andrea Thamm[2]
[1]*Institut für Theoretische Physik, Universität Heidelberg, Philosophenweg 16, 69120 Heidelberg, Germany*
[2]*PRISMA Cluster of Excellence & MITP, Johannes Gutenberg University, 55099 Mainz, Germany*
[3]*Department of Physics & LEPP, Cornell University, Ithaca, New York 14853, USA*

(Received 4 July 2016; revised manuscript received 24 August 2016; published 28 October 2016)

Scalar particles S that are singlets under the standard model gauge group are generic features of many models of fundamental physics, in particular, as possible mediators to a hidden sector. We show that the decay $S \to Zh$ provides a powerful probe of the CP nature of the scalar, because it is allowed only if S has CP-odd interactions. We perform a model-independent analysis of this decay using an effective Lagrangian and compute the relevant Wilson coefficients arising from integrating out heavy fermions to one-loop order.

DOI: 10.1103/PhysRevLett.117.181801

Introduction.—Pseudoscalar singlets play an important role in various extensions of the standard model (SM). They appear, e.g., as mediators to a dark sector or in solutions to the strong CP problem. Searches at the LHC focus on the model-specific signals of these new states, which often do not reveal their pseudoscalar nature—the phantom digamma excess seen in the first 13 TeV data [1,2] could have been an example of such a signal. Identifying the CP properties of such a new state will be one of the top priorities if a signal is seen in future data.

Let us consider a new spin-0 particle S, which is a gauge singlet under the SM gauge group. Assuming its mass is much larger than the electroweak scale, its interactions can be described in terms of local operators in the unbroken phase of the electroweak gauge symmetry. At the renormalizable level, the only interactions of S with SM particles arise from the Higgs portals

$$\mathcal{L}_{\text{portal}} = -\lambda_1 S\phi^\dagger\phi - \frac{\lambda_2}{2} S^2\phi^\dagger\phi, \tag{1}$$

where ϕ is the Higgs doublet. The first term gives rise to a mixing between S and the Higgs boson, with a mixing angle $\alpha \sim v\lambda_1/m_S^2$. The coupling λ_1 is naturally of the order of the UV cutoff of the theory, but at least of order m_S, and hence one expects $\alpha > v/m_S$. However, this mixing affects the phenomenology of Higgs decay rates, and so in practice α must be small. The example of the elusive 750 GeV diphoton resonance [1,2] has demonstrated that tight bounds on α can also be derived from the decays $S \to ZZ$, WW, $t\bar{t}$, hh [3,4]. The portal coupling λ_2, on the other hand, does not give rise to dangerous effects.

It is therefore a challenge to model building to find ways of suppressing the coupling λ_1, either by means of a symmetry or dynamically. A discrete Z_2 symmetry under which S changes sign would enforce $\lambda_1 = 0$. If the ultraviolet theory is (at least approximately) CP invariant, then neutral particles can be classified as CP eigenstates. If S is a CP-odd pseudoscalar ($J^{PC} = 0^{-+}$), λ_1 must be 0. A nice example of a dynamical suppression is provided by models in which S is identified with the lowest mode of a Z_2-odd bulk scalar in a warped extra dimension [3,5]. When the Higgs sector is localized on the IR brane, its coupling to S is either suppressed by a small wave-function overlap or by a loop factor. Here we entertain the possibility of eliminating the portal coupling λ_1 by supposing that S is a CP-odd pseudoscalar, e.g., an axionlike particle.

Measurements of angular distributions in $S \to ZZ \to 4l$ or $S \to Z\gamma \to 4l$ decays have been considered as a way of probing the spin and CP properties of a new resonance [6,7], in analogy with the corresponding measurements in Higgs decays [8]. However, the rates for these decays are likely to be quite small, since a gauge-singlet S has no renormalizable couplings to gauge bosons. Hence it may require very large statistics to perform these analyses. In this Letter we propose the decay $S \to Zh$, which is strictly forbidden for a CP-even scalar, as a novel and independent way to test the spin and CP quantum numbers of a new particle S. The very existence of this decay would constitute a smoking-gun signal for a pseudoscalar nature of S (or for significant CP-odd couplings, in case S is a state with mixed CP quantum numbers), without the need to analyze angular distributions. The observation of this decay would also exclude a spin-2 explanation of a hypothetical new resonance [9]. To the best of our knowledge this signature has not been studied in the literature. Established experimental searches in the context of two-Higgs-doublet models can be adapted for the proposed search. The most promising decay mode is $S \to Zh \to l^+l^-b\bar{b}$ [10].

Effective Lagrangian analysis.—At the level of dimension-5 operators, the most general couplings of a CP-odd scalar to gauge bosons read

$$\mathcal{L}_{\text{eff}}^{\text{gauge}} = \frac{\tilde{c}_{gg}}{M}\frac{\alpha_s}{4\pi} S G_{\mu\nu}^a \tilde{G}^{\mu\nu,a} + \cdots, \tag{2}$$

where M denotes the new-physics scale, and the dots represent analogous couplings to the $SU(2)_L$ and $U(1)_Y$

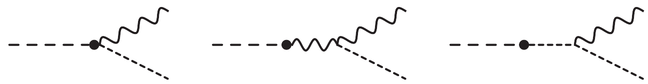

FIG. 1. Tree-level diagrams representing the contribution of the operator in (3) to $S \to Zh$ decay. The internal dashed line in the third graph represents the Goldstone boson φ_3.

gauge bosons. Via this operator the resonance S can be produced in gluon fusion at the LHC. The most general dimension-5 couplings of S to fermions have the same form as the SM Yukawa interactions times S/M, and with the Yukawa matrices replaced by some new matrices. In any realistic model these couplings must have a hierarchical structure in the mass basis in order to be consistent with the strong constraints from flavor physics [11]. It is thus reasonable to assume that the dominant couplings are those to the top quarks; see (5) below.

When using an effective Lagrangian to describe the production and decays of the resonance S one should keep in mind that in many new-physics scenarios the masses of the heavy particles which are integrated out are in the TeV range. When there is no significant mass gap between S and the new sector, contributions from operators with dimension $D \geq 6$ are not expected to be strongly suppressed. Some of these operators can induce new structures not present at dimension-5 level.

$D = 5$ operator analysis of $S \to Zh$ decay: The decay $S \to Zh$ has been studied in the context of two-Higgs-doublet models, where it arises at the renormalizable level via the kinetic terms [12,13]. However, this requires the pseudoscalar S to be light (since the effect vanishes in the decoupling limit) and carry electroweak quantum numbers. In this case the existence of CP-odd couplings of the heavy scalar bosons can be related to three $U(2)$ invariants of the scalar potential [14]. For the case of a gauge-singlet scalar considered here no such invariants exist. Moreover, the effective Lagrangian up to dimension 5 does not contain any polynomial operator that could mediate the decay $S \to Zh$ at tree level. The obvious candidate

$$(\partial^\mu S)(\phi^\dagger iD_\mu\phi + \text{H.c.}) \to -\frac{g}{2c_w}(\partial^\mu S)Z_\mu(v+h)^2, \qquad (3)$$

where $c_w \equiv \cos\theta_w$ and the second expression holds in unitary gauge, can be reduced to operators containing fermionic currents using the equations of motion. This follows from the partial conservation of the Higgs current

$$\partial^\mu(\phi^\dagger iD_\mu\phi + \text{H.c.}) \to -\left(1+\frac{h}{v}\right)\sum_f 2T_3^f m_f \bar{f} i\gamma_5 f, \qquad (4)$$

where T_3^f is the third component of weak isospin. The resulting operators do not give rise to a tree-level $S \to Zh$ matrix element. Indeed, adding up the diagrams shown in Fig. 1 one finds that the tree-level $S \to Zh$ matrix element of the operator in (3) vanishes identically, and the same is true for the $S \to Zhh$ matrix element.

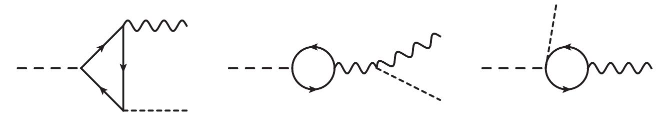

FIG. 2. Top-loop contributions to $S \to Zh$ decay. We omit a mirror copy of the first graph with a different orientation of the fermion loop and diagrams involving Goldstone bosons.

At one-loop order, the $S \to Zh$ decay amplitude receives a contribution from an operator containing quark fields, and since the Higgs boson couples proportional to the quark mass it suffices to consider the term involving the top quark. The relevant Lagrangian is

$$\mathcal{L}_{\text{eff}}^{D=5} = -\tilde{c}_{tt}\frac{y_t}{M}S(i\bar{Q}_L\tilde{\phi}t_R + \text{H.c.}), \qquad (5)$$

where Q_L is the third-generation left-handed quark doublet and $\tilde{\phi} = \epsilon\phi^*$. The one-loop Feynman diagrams contributing to the decay $S \to Zh$ are shown in Fig. 2. Analogous diagrams involving electroweak gauge bosons in the loop vanish, since it is impossible to saturate the Lorentz indices of the $\epsilon^{\mu\nu\alpha\beta}$ tensor associated with the dual field strength in CP-odd interactions such as (2). We have evaluated the diagrams in Fig. 2 in a general R_ξ gauge. The resulting decay amplitude is

$$i\mathcal{A}(S \to Zh) = -\frac{2m_Z\epsilon_Z^* \times p_h}{M}C_5^{\text{top}},$$
$$\text{with} \quad C_5^{\text{top}} = -\frac{N_c y_t^2}{8\pi^2}T_3^t\tilde{c}_{tt}F, \qquad (6)$$

where $T_3^t = \frac{1}{2}$. The Z boson is longitudinally polarized, and hence the structure $2m_Z\epsilon_Z^* \cdot p_h \approx 2p_Z \cdot p_h \approx m_S^2$ is proportional to the mass squared of the heavy particle. The quantity F denotes the parameter integral

$$F = \int_0^1 d[xyz]\frac{2m_t^2 - xm_h^2 - zm_Z^2}{m_t^2 - xzm_S^2 - xym_h^2 - yzm_Z^2 - i0}, \qquad (7)$$

with $d[xyz] \equiv dxdydz\delta(1-x-y-z)$. The factor $y_t^2 = 2m_t^2/v^2$ in (6) ensures that analogous contributions from light fermions in the loop are negligible. Evaluating the integral with $m_t \equiv m_t(m_S)$ and with the physical Higgs and Z-boson masses gives $F \approx -0.010 + 0.673i$ for $m_S = 750$ GeV and $F \approx -0.092 + 0.230i$ for $m_S = 1.5$ TeV, where here and below we pick two representative values for the mass of the pseudoscalar resonance. For $m_S^2 \gg m_t^2$, the function F is formally suppressed by a factor m_t^2/m_S^2, but its imaginary part is numerically enhanced. From the amplitude (6) we obtain the decay rate

$$\Gamma(S \to Zh)_{D=5} = \frac{m_S^3}{16\pi M^2}|C_5^{\text{top}}|^2\lambda^{3/2}(1, x_h, x_Z), \qquad (8)$$

where $x_i = m_i^2/m_S^2$ and $\lambda(x,y,z) = (x-y-z)^2 - 4yz$. We find $\Gamma(S \to Zh)_{D=5} \approx 0.6\ \text{MeV}\tilde{c}_{tt}^2(\text{TeV}/M)^2$ in both cases.

Assuming that the dominant contribution to the $S \to Zh$ decay amplitude indeed arises at dimension 5, one can derive the model-independent relation

$$\frac{\Gamma(S \to Zh)_{D=5}}{\Gamma(S \to t\bar{t})} = \frac{3y_t^2}{16\pi^2}\left(\frac{m_S}{4\pi v}\right)^2 |F|^2 \frac{\lambda^{3/2}(1, x_h, x_Z)}{\sqrt{1-4x_t}}. \quad (9)$$

This ratio evaluates to 3.6×10^{-4} for $m_S = 750$ GeV and 1.8×10^{-4} for $m_S = 1.5$ TeV. The present experimental upper bounds on the corresponding $S \to t\bar{t}$ rates of about 0.7 pb and 65 fb at $\sqrt{s} = 8$ TeV [15] yield $\sigma(pp \to S_{750} \to t\bar{t}) < 3.2$ pb and $\sigma(pp \to S_{1500} \to t\bar{t}) < 0.6$ pb at $\sqrt{s} = 13$ TeV under the assumption of gluon-initiated production. Relation (9) then implies the bounds $\sigma(pp \to S_{750} \to Zh)_{D=5} < 1.1$ fb and $\sigma(pp \to S_{1500} \to Zh)_{D=5} < 0.1$ fb, which are 2 orders of magnitude below the direct experimental upper limits $\sigma(pp \to S_{750} \to Zh) < 123$ fb and $\sigma(pp \to S_{1500} \to Zh) < 40$ fb at $\sqrt{s} = 13$ TeV [10]. Note that the former bounds do not apply if $m_S < 2m_t$.

$D = 7$ operator analysis of $S \to Zh$ decay: The dominance of the loop-induced dimension-5 contribution to the $S \to Zh$ decay rate is far from guaranteed. This contribution can be very small if the CP-odd coupling $\tilde{c}_{tt}$ of S to top quarks is suppressed. Also, as we have seen, the one-loop matrix element in (6) is suppressed by a factor m_t^2/m_S^2. If m_S is not much smaller than the new-physics scale M, the loop contributions arising at dimension 7 can give rise to similar effects. Moreover, at dimension 7 there exists a unique operator giving rise to a tree-level contribution to the $S \to Zh$ amplitude. It reads

$$\begin{aligned} O_7 &= (\partial^\mu S)(\phi^\dagger iD_\mu \phi + \text{H.c.})\phi^\dagger\phi \\ &\triangleq -S(\phi^\dagger iD_\mu\phi + \text{H.c.})\partial^\mu(\phi^\dagger\phi) \\ &\to \frac{g}{2c_w} SZ_\mu (v+h)^3 \partial^\mu h, \end{aligned} \quad (10)$$

where in the second step we have used an integration by parts and the equations of motion for the Higgs field, neglecting the fermionic terms in (4), which do not contribute to $S \to Zh$ decay at tree level. The expression in the third line, valid in unitary gauge, gives rise to nonvanishing $S \to Zh$ and $S \to Zhh$ matrix elements.

At one-loop order there exist several dimension-7 operators contributing to the decay $S \to Zh$. Those that mix with O_7 under renormalization are

$$\begin{aligned} \mathcal{L}_{\text{eff}}^{D=7} &= \frac{C_7}{M^3} O_7 + \frac{c_6^t}{M^2} \bar{t}_R \tilde{\phi}^\dagger i\not{D}\tilde{\phi} t_R \\ &+ \frac{c_{7a}^t}{M^3}(iS\bar{Q}_L i\not{D} i\not{D}\tilde{\phi} t_R + \text{H.c.}) \\ &+ \frac{c_{7b}^t}{M^3}(\partial^\mu S)\bar{t}_R \tilde{\phi}^\dagger \gamma^\mu \tilde{\phi} t_R + \cdots, \end{aligned} \quad (11)$$

plus analogous operators containing the right-handed bottom quark. The dimension-6 operator proportional to c_6^t contributes in conjunction with the operator in (5) to give a contribution of order $1/M^3$.

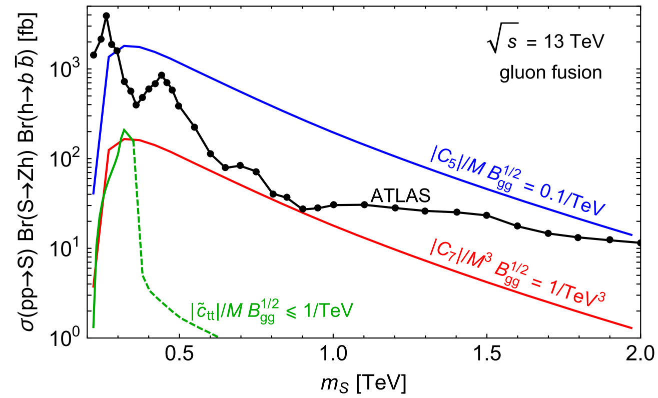

FIG. 3. Predictions for the $pp \to S \to Zh \to Zb\bar{b}$ signal rate vs m_S, compared with the ATLAS upper bounds [10]. The red line shows the contribution from C_7 evaluated with $B_{gg}^{1/2}|C_7|/M^3 = 1/\text{TeV}^3$, while the blue line shows a generic dimension-5 contribution with $B_{gg}^{1/2}|C_5|/M = 0.1/\text{TeV}$ (see Section "Nonpolynomial operators"), where $B_{gg} \equiv \text{Br}(S \to gg)$. The green line shows the contribution from C_5^{top} for $B_{gg}^{1/2}|\tilde{c}_{tt}|/M = 1/\text{TeV}$, while the dashed green line incorporates the upper bound on $|\tilde{c}_{tt}|$ implied by the ATLAS limits on the $pp \to S \to t\bar{t}$ rate [15].

Let us focus on the potentially dominant tree-level contribution from O_7, which yields the decay rate

$$\Gamma(S \to Zh) \approx \frac{m_S^3}{16\pi M^2}\left|C_5^{\text{top}} + \frac{v^2}{2M^2}C_7\right|^2 \lambda^{3/2}(1, x_h, x_Z). \quad (12)$$

With $C_7 = 1$ and $M = 1$ TeV this partial width is about 7 MeV for $m_S = 750$ GeV and 60 MeV for $m_S = 1.5$ TeV. The contribution from C_5^{top} can be safely neglected in this case, except in the kinematic region where $m_S < 2m_t$. Figure 3 shows our results for the $pp \to S \to Zh \to Zb\bar{b}$ signal rate under the assumption that S is produced in gluon fusion [3] and that a single Wilson coefficient gives the dominant contribution to the $S \to Zh$ rate. We fix the products $B(S \to gg)^{1/2}|C_7|/M^3$ etc. to the values shown in the plot. The rate scales with the squares of these combinations. These results show that $S \to Zh$ rates close to the present experimental bounds are possible for reasonable parameter values, provided that the $S \to gg$ branching ratio is not too small.

Nonpolynomial operators: It is interesting to consider the hypothetical limit $m_t \gg m_S$ in (7). Then the parameter integral yields $F = 1 + \mathcal{O}(m_S^2/m_t^2)$. The fermion is a very heavy particle, which can be integrated out from the low-energy theory. The contribution (6) then corresponds to a one-loop matching contribution to the Wilson coefficient of a local dimension-5 operator with a tree-level $S \to Zh$ matrix element. Our operator analysis in Section "$D = 5$ operator analysis of $S \to Zh$ decay" did not reveal the existence of such an operator. However, in extensions of the

SM containing heavy particles whose masses arise (or receive their dominant contributions) from electroweak symmetry breaking, operators with a nonpolynomial dependence on the Higgs field can arise [16]. The nonpolynomial structure appears because the particle integrated out (the hypothetical heavy fermion) receives its mass from electroweak symmetry breaking, so it is heavy only in the broken phase of the theory. In our case, the relevant operator reads

$$\begin{aligned} O_5 &= (\partial^\mu S)(\phi^\dagger iD_\mu\phi + \text{H.c.})\ln\frac{\phi^\dagger\phi}{\mu^2} \\ &\triangleq -S(\phi^\dagger iD_\mu\phi + \text{H.c.})\frac{\partial^\mu(\phi^\dagger\phi)}{\phi^\dagger\phi}, \end{aligned} \tag{13}$$

where in the second step we have again used an integration by parts and neglected fermionic currents. The latter expression has a one-to-one map onto the structure of the parameter integral (7).

Consider, as an illustration, a sequential fourth generation of heavy leptons, and assume that the heavy charged state L has a mass $m_L \gtrsim m_S/2$ and a coupling $\tilde{c}_{LL}$ to the pseudoscalar resonance defined in analogy to (5). Integrating out this heavy lepton generates the contribution

$$C_5 = \frac{y_L^2\tilde{c}_{LL}}{16\pi^2} = \frac{m_L^2\tilde{c}_{LL}}{8\pi^2v^2} \gtrsim \frac{m_S^2\tilde{c}_{LL}}{32\pi^2v^2} \tag{14}$$

to the Wilson coefficient C_5/M of the operator O_5. Using this expression instead of C_5^{top} in (8), we obtain the upper bounds $|\tilde{c}_{LL}| < 1.3(M/\text{TeV})$ for $m_S = 750$ GeV and $|\tilde{c}_{LL}| < 0.6(M/\text{TeV})$ for $m_S = 1.5$ TeV. In such a model it would be natural to obtain $S \to Zh$ decay rates close to the present experimental upper bounds, see Fig. 3.

The effective Lagrangian $\mathcal{L}_{\text{eff}} = (C_5/M)O_5$ yields a loop correction to the T parameter given by $\alpha(m_Z)T = -\Pi_{ZZ}(0)/m_Z^2 \approx C_5^2/(4\pi)^2$. Electroweak precision measurements then imply $|C_5| < 0.66$ at 95% confidence level [17]. This constraint is much weaker than the bounds derived from $S \to Zh$ decay.

The operator O_5 and analogous nonpolynomial operators of higher dimension are absent in models where the new heavy particles have masses not related to the electroweak scale. We now study such a model in detail.

Heavy vectorlike fermions.—It is instructive to consider a concrete new-physics model, which generates the effective interactions of the scalar resonance with SM particles via loop diagrams involving heavy vectorlike fermions that are mixed with the SM fermions. Such a scenario is realized, e.g., in models of partial compositeness or warped extra dimension [18–20]. We consider an $SU(2)_L$ doublet $\psi = (TB)^T$ of vectorlike quarks with hypercharge $Y_\psi = \frac{1}{6}$, which mixes with the third-generation quark doublet of the SM. The most general Lagrangian reads

$$\begin{aligned} \mathcal{L} &= \bar{\psi}(i\not{D} - M)\psi + \bar{Q}_L i\not{D} Q_L + \bar{t}_R i\not{D} t_R + \bar{b}_R i\not{D} b_R \\ &\quad - y_t(\bar{Q}_L\tilde{\phi}t_R + \text{H.c.}) - (g_t\bar{\psi}\,\tilde{\phi}\,t_R + g_b\bar{\psi}\phi b_R + \text{H.c.}) \\ &\quad - c_1 S\bar{\psi}i\gamma_5\psi - ic_2 S(\bar{Q}_L\psi - \bar{\psi}Q_L), \end{aligned} \tag{15}$$

where we neglect the small Yukawa coupling $|y_b| \ll 1$ of the bottom quark. The terms in the last line contain the couplings to the pseudoscalar resonance S. The mass mixing induced by the couplings g_i leads to modifications of the masses and Yukawa couplings of the SM top and bottom quarks by small amounts of order $g_i^2v^2/M^2$. Likewise, the masses of the heavy T and B quarks are split by a small amount $M_T - M_B \approx (g_t^2 - g_b^2)v^2/(4M)$.

Integrating out the heavy fermion doublet at tree level, by solving its equations of motion, we generate the operators in the effective Lagrangians (5) and (11) with coefficients $\tilde{c}_{tt} = -c_2g_t/y_t$ and (for $f = t, b$)

$$c_6^f = g_f^2, \qquad c_{7a}^f = c_2g_f, \qquad c_{7b}^f = c_1g_f^2. \tag{16}$$

The coefficient c_6^b is constrained by precision measurements of the Z-boson couplings to fermions performed at LEP and SLD. A recent global analysis finds [21]

$$c_6^b = g_b^2 = (0.76 \pm 0.27)\left(\frac{M}{\text{TeV}}\right)^2, \tag{17}$$

where the pull away from 0 is largely driven by the b-quark forward-backward asymmetry A_b^{FB}, whose experimental value is about 2.8σ smaller than the SM prediction [22]. Our model can resolve this anomaly in a natural way. It is likely that the coupling g_t is at least as large as g_b, perhaps even significantly larger. In our model the relation $\tilde{c}_{bb}/\tilde{c}_{tt} = (g_b/g_t)(m_t/m_b)$ holds, and hence the coupling of the resonance S to bottom quarks defined in analogy with (5) can be rather large.

The coefficient C_7 in (11) is induced at one-loop order by diagrams such as those shown in Fig. 2, where now both heavy and light quarks can propagate in the loops. In order to calculate C_7 a proper matching onto the low-energy theory must be performed. We obtain

$$\begin{aligned} \frac{v^2}{2}C_7 &= c_1\sum_{f=t,b}\frac{N_cg_f^2}{16\pi^2}\left\{2T_3^f\left[m_f^2\left(L - \frac{3}{2}\right)\right.\right. \\ &\quad \left.\left. - \frac{m_h^2}{12} + \frac{m_Z^2}{36} + \frac{g_f^2v^2}{4}\right] - \frac{2}{3}Q_fs_w^2m_Z^2\left(L - \frac{3}{2}\right)\right\} \\ &\quad + \tilde{c}_{tt}\frac{N_cy_t^2}{16\pi^2}\left\{2T_3^t\left[3m_t^2\left(L - \frac{3}{2}\right)\right.\right. \\ &\quad - \frac{m_h^2}{2}\left(L - \frac{7}{6}\right) - \frac{m_Z^2}{6}\left(L + \frac{19}{6}\right) \\ &\quad \left.\left. - g_t^2v^2\left(L - \frac{9}{4}\right)\right] + Q_ts_w^2m_Z^2\right\}, \end{aligned} \tag{18}$$

where $L = \ln(M^2/\mu^2)$. Note the absence of terms proportional to m_S^2 on the right-hand side of this expression, which is a consequence of the fact that there is no corresponding dimension-7 operator. There is a nontrivial operator mixing, such that the scale dependence of the coefficient C_7 cancels against the scale dependence of the one-loop matrix elements of the fermionic operators in the effective Lagrangian (11); see [23] for details. To estimate the dimension-7 contribution we set $\mu = m_Z$ in (18) and neglect the fermion-loop contributions in the low-energy theory. All large logarithms $L \approx 4.8$ are included in the Wilson coefficient C_7, for which we obtain, assuming $M \approx 1$ TeV in the argument of the logarithms, the expression

$$C_7 \approx [c_1(5.30g_t^2 + 0.95g_t^4 + 0.16g_b^2 - 0.95g_b^4) + \tilde{c}_{tt}(10.18 - 6.90g_t^2)] \times 10^{-2}. \quad (19)$$

For natural values of the couplings this coefficient can be rather large. For example, with $g_t = 2$ and $g_b = 0.87$ set by (17) we get $C_7 \approx (0.36c_1 - 0.17\tilde{c}_{tt})$. For a 750 GeV resonance produced in gluon fusion (and dominantly decaying to dijets) the production rate

$$\sigma(pp \to S)\mathrm{Br}(S \to Zh) \approx 70\ \mathrm{fb}\left(\frac{\mathrm{TeV}}{M}\right)^6 C_7^2 \quad (20)$$

can be sizeable and, especially for $M < 1$ TeV, can even come close to the current upper bound of 123 fb [10]. For $m_S = 1.5$ TeV the rate is smaller by about a factor 10.

Conclusions.—We have presented the first detailed analysis of the decay $S \to Zh$ of a gauge-singlet, heavy spin-0 particle S and pointed out that this process is allowed only if S has CP-odd interactions. Such a pseudoscalar boson arises in many well-motivated extensions of the SM, including models containing Higgs-portal mediators to a hidden sector and scenarios addressing the strong CP problem. Alternative ways to determine the CP nature of a new boson rely on high statistics to perform analyses of angular distributions, whereas the mere observation of the decay $S \to Zh$ proposed here would establish the presence of a CP-odd coupling.

Using a model-independent analysis based on an effective Lagrangian, we have shown that the decay amplitude receives fermion-loop contributions starting at dimension 5, while tree-level contributions can first arise at dimension 7. In new-physics models containing heavy particles whose masses arise from electroweak symmetry breaking there also exists a nonpolynomial dimension-5 operator with a tree-level $S \to Zh$ matrix element.

We have derived explicit expressions for the relevant Wilson coefficients at one-loop order in a model containing heavy vectorlike fermions with CP-odd couplings to S, finding that appreciable $pp \to S \to Zh$ production rates, even close to the present experimental bounds, can be obtained for reasonable values of parameters. This motivates a vigorous experimental program to search for $S \to Zh$ decays in the high-luminosity LHC run.

M. B. acknowledges the support of the Alexander von Humboldt Foundation. The work of M. N. and A. T. is supported by the Advanced Grant No. EFT4LHC of the European Research Council (ERC), the DFG Cluster of Excellence PRISMA (Grant No. EXC 1098), and Grant No. 05H12UME of the German Federal Ministry for Education and Research (BMBF).

[1] M. Aaboud *et al.* (ATLAS Collaboration), J. High Energy Phys. 09 (2016) 001.
[2] V. Khachatryan *et al.* (CMS Collaboration), Phys. Rev. Lett. **117**, 051802 (2016).
[3] M. Bauer, C. Hörner, and M. Neubert, J. High Energy Phys. 07 (2016) 094.
[4] S. Dawson and I. M. Lewis, arXiv:1605.04944.
[5] C. Csaki and L. Randall, J. High Energy Phys. 07 (2016) 061.
[6] M. Chala, C. Grojean, M. Riembau, and T. Vantalon, Phys. Lett. B **760**, 220 (2016).
[7] R. Franceschini, G. F. Giudice, J. F. Kamenik, M. McCullough, F. Riva, A. Strumia, and R. Torre, J. High Energy Phys. 07 (2016) 150.
[8] A. Soni and R. M. Xu, Phys. Rev. D **48**, 5259 (1993).
[9] H. M. Lee, D. Kim, K. Kong, and S. C. Park, J. High Energy Phys. 11 (2015) 150.
[10] ATLAS Collaboration, Report No. ATLAS-CONF-2016-015.
[11] F. Goertz, J. F. Kamenik, A. Katz, and M. Nardecchia, J. High Energy Phys. 05 (2016) 187.
[12] H. Baer, C. Kao, and X. Tata, Phys. Lett. B **303**, 284 (1993).
[13] D. Kominis, Nucl. Phys. **B427**, 575 (1994).
[14] J. F. Gunion and H. E. Haber, Phys. Rev. D **72**, 095002 (2005).
[15] G. Aad *et al.* (ATLAS Collaboration), J. High Energy Phys. 08 (2015) 148.
[16] A. Pierce, J. Thaler, and L. T. Wang, J. High Energy Phys. 05 (2007) 070.
[17] M. Baak, J. Cúth, J. Haller, A. Hoecker, R. Kogler, K. Mönig, M. Schott, and J. Stelzer (Gfitter Group Collaboration), Eur. Phys. J. C **74**, 3046 (2014).
[18] D. B. Kaplan, Nucl. Phys. **B365**, 259 (1991).
[19] Y. Grossman and M. Neubert, Phys. Lett. B **474**, 361 (2000).
[20] T. Gherghetta and A. Pomarol, Nucl. Phys. **B586**, 141 (2000).
[21] A. Efrati, A. Falkowski, and Y. Soreq, J. High Energy Phys. 07 (2015) 018.
[22] S. Schael *et al.* (ALEPH, DELPHI, L3, OPAL, SLD, LEP Electroweak Working Group, SLD Electroweak Group, and SLD Heavy Flavour Group Collaborations), Phys. Rep. **427**, 257 (2006).
[23] M. Bauer, M. Neubert, and A. Thamm, arXiv:1607.01016.