Polarization effects in the Higgs boson decay to $\gamma Z$ and test of $CP$ and $CPT$ symmetries

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Polarization characteristics of $\gamma$ and $\gamma Z$ states in the Higgs boson decays $h \to \gamma \gamma$ and $h \to \gamma Z$ are discussed. Based on effective Lagrangian, describing $h \gamma \gamma$ and $h \gamma Z$ interactions with $CP$-even and $CP$-odd parts, we calculate polarization parameters $\xi_1$, $\xi_2$, $\xi_3$. A nonzero value of the photon circular polarization, defined by parameter $\xi_3$, arises due to presence of both parts in effective Lagrangian and its non-Hermiticity. The circular polarization is proportional to the forward-backward asymmetry of fermions in the decay $h \to \gamma Z \to f \bar{f}$. Measurement of this observable would allow one to search for deviation from the standard model and possible violation of $CPT$ symmetry. We discuss also a possibility to measure parameters $\xi_1$, $\xi_3$, describing correlation of linear polarizations of photon and $Z$ boson, in the decay $h \to \gamma^* Z \to \ell^+ \ell^-$ via distribution over the azimuthal angle between the decay planes of $\gamma^* \to \ell^+ \ell^-$ and $Z \to \bar{f} f$. Deviation of the measured value of $\xi_1$ from zero will indicate $CP$ violation in the Higgs sector.

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I. INTRODUCTION

The ATLAS and CMS collaborations at the LHC have recently observed [1, 2] a boson $h$ with mass around 126 GeV with statistical significance of about five standard deviations. The experimental evidence of this new particle is the strongest in the two-photon and four-lepton final channels, where the detectors give the best mass resolution.

Although the decay pattern of $h$ is mainly consistent with the predictions of the standard model (SM), the clarification of the nature of this particle still needs more data and time. The spin of this boson is known to be zero or two, while the $CP$ properties are not yet ascertained.

Recent data are more consistent with the pure scalar boson hypothesis than the pure pseudoscalar one [3]. Though in the SM the Higgs boson has $J^{PC} = 0^{++}$, there are many extensions of the SM with a more complicated Higgs sector, in which some of the Higgs bosons may not have definite $CP$ parity [4].

This aspect of the Higgs study is also related to the origin of the $CP$ violation. In the SM the source of the $CP$ violation is the complex irreducible phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix [5], however this effect is not sufficient to explain the observed matter-antimatter asymmetry in the Universe [6]. There may be other mechanisms of the $CP$ violation beyond the CKM matrix, for example, in the Higgs sector. From this point of view, the elucidation of the $CP$ properties of the observed $h$ boson would be an important step towards clarification of the mechanisms giving rise to the masses of particles, their mixing and $CP$ violation.

Recently the $CP$ properties of the Higgs boson in the two-photon decay channel $h \to \gamma \gamma$ have been addressed in Ref. [9]. In this channel the branching fraction, measured by the ATLAS collaboration, is larger than the value predicted in the SM by a factor of $1.60 \pm 0.30$ for $m_h = 125.2 \pm 0.26$ (stat)$^{+0.5}_{-0.6}$ (syst) GeV [10], while the CMS collaboration obtained for this factor $0.77 \pm 0.27$ for $m_h = 125.7 \pm 0.3$ (stat) $\pm 0.3$ (syst) GeV [11]. The author of [9], in framework of a model with vectorlike fermions, showed that the $CP$ violation in the $h \to \gamma \gamma$ decay results in the dependence of the differential decay rate on the angle between linear polarization vectors of the photons. Experimentally, this angular distribution can be measured after both photons are converted into the $e^+, e^-$ pairs via the azimuthal angle distribution between the planes spanned by the two $e^+$, $e^-$ pairs. In Ref. [12] a model-independent analysis of the $CP$ violation effects in the Higgs boson into a pair of the gauge bosons $W^+$, $W^-$ or $Z$, $Z$ has been presented. The author has studied the angular distributions of the fermions $f = \ell, q$ in the cascade processes $h \to V_1 V_2 \to (f_1 f_2) (\bar{f}_3 \bar{f}_4)$ and analyzed possibilities of observation of the $CP$ violation in these decays to various final lepton and quark pairs.

In the present paper we would like to address the decay of the Higgs boson to the photon and $Z$ boson, $h \to \gamma Z$, pointing out to a possibility of studying in this decay not only the $CP$ properties of the newly discovered boson, but also the validity of the $CPT$ symmetry. In this connection one can recall Ref. [13] in which the author showed that an observation of the circular polarization of the photon in the neutral pion decay $\pi^0 \to \gamma \gamma$ (or $\eta \to \gamma \gamma$) would signal violation of the $CPT$ symmetry.

Indeed, the product $s \bar{k}$ (where $s$ is the photon spin and $\bar{k}$ is its momentum) is $P$ odd and $T$ even. Such a correlation in the $\pi^0$ decay arises due to interference of the two terms in the interaction Lagrangian: a scalar $\epsilon \pi^0 \epsilon^{\mu \nu \sigma} F_{\mu \nu} F_{\rho \sigma}$ and a pseudoscalar $c \pi^0 F_{\mu \nu} F^{\mu \nu}$, with $c$ and $\epsilon$ being con-
plings constants and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. The analysis of \cite{13} demonstrated that a nonzero value of $\delta k$ correlation may appear due to a non-Hermiticity of the three-level amplitude, i.e. $\Im \overline{\epsilon} \neq 0$ or/and $\Im m \neq 0$, and/or higher-order loop corrections to the amplitude inducing imaginary part of $\overline{\epsilon}$.

Note that such a correlation in the Higgs boson decay to two transversally polarized $Z$ bosons in connection with possible violation of CPT symmetry has been discussed in \cite{12}.

Generally, similar arguments can be applied to the two-photon decay of the Higgs boson with an analogous conclusion. However measurement of the photon circular polarization in the $h \to \gamma \gamma$ decay is a rather difficult task. In the present paper we suggest to study CP and possible CPT violation in the decay

$$h \to \gamma Z \to \gamma f \bar{f},$$

with $f = \ell, q$. It turns out that the decay distribution over the angle $\theta$ between the momentum of the fermion $f$ (in the rest frame of the $Z$) and momentum of the $Z$ (in the rest frame of the $h$) gives information on the photon circular polarization. Namely, a nonzero photon circular polarization induces a term $\sim \cos \theta$ in this distribution which can be measured through the forward-backward asymmetry $A_{\text{FB}}$.

In the SM the $h \to \gamma Z$ decay amplitude in the lowest order is determined by the loop contributions \cite{14,15}, which have a small but nonzero imaginary part arising due to rescattering effects $h \to f \bar{f} \to \gamma Z$ for the fermions $f$ with masses $m_f \leq m_h / 2$. The corresponding effective Lagrangian $\mathcal{L}_{\text{eff}}^{h\gamma Z}$, describing interaction of $h, \gamma$ and $Z$, is thus non-Hermitian. Non-Hermiticity of effective Lagrangian leads to a nonzero value of the net photon helicity once we assume a mixture of CP violating term in $\mathcal{L}_{\text{eff}}^{h\gamma Z}$. Note that in the SM and theories beyond the SM which are CPT symmetric, there are no sources of non-Hermiticity of $\mathcal{L}_{\text{eff}}^{h\gamma Z}$ apart from rescattering effects.

The CPT theorem is one of the most profound results of quantum field theory \cite{12}. It is a consequence of Lorentz invariance, locality, connection between spin and statistics, and a Hermitian Hamiltonian. However there are many extensions of the SM in which CPT violation appears due to nonlocality in the string theory, or violation of Lorentz symmetry in the extra dimensional models (see, for example, \cite{17}). One can also mention possible deviations from the standard quantum mechanical evolution of states in some models of quantum gravity, and the corresponding breakdown of the CPT symmetry is investigated in the neutral-meson system, where novel CPT-violating observables for the $\phi$-factories and $B$-factories are proposed \cite{18}. The CPT violating effects in some of these underlying theories, in principle, can be additional sources of non-Hermiticity of effective Lagrangian $\mathcal{L}_{\text{eff}}^{h\gamma Z}$ and hence contribute to photon circular polarization.

As for experimental results on the SM Higgs boson decay to the $Z$ boson and photon, we mention recent ATLAS and CMS results \cite{19,20}. The Higgs production cross section times the $h \to \gamma Z$ branching fraction limits are about an order of magnitude larger than the SM expectation for $m_h = 125$ GeV.

The paper is organized as follows. In Sec. II effective Lagrangian for $h \gamma \gamma$ and $h \gamma Z$ interactions and coupling constants in the SM and some its extensions are considered. In Sec. III amplitudes and polarization parameters for the decays $h \to \gamma \gamma$ and $h \to \gamma Z$ are specified. Distribution of the $h \to \gamma Z \to \gamma f \bar{f}$ decay in the polar angle, and distribution of the $h \to \gamma^* Z \to \ell^+ \ell^- Z$ decay (with $Z \to f f$ on mass shell) in the azimuthal angle are obtained. In Sec. IV results of calculation and discussion are presented. In Sec. V we draw conclusions.

II. FORMALISM

The effective Lagrangian for the $h \gamma \gamma$ and $h \gamma Z$ interactions can be written as

$$\mathcal{L}_{\text{eff}}^{h\gamma\gamma} = \frac{\epsilon^2}{32 \pi^2 v} \left( c_\gamma F_{\mu\nu} F^{\mu\nu} h - \overline{\epsilon}_\gamma F_{\mu\nu} \overline{F}^{\mu\nu} h \right),$$

$$\mathcal{L}_{\text{eff}}^{h\gamma Z} = \frac{\epsilon g}{16 \pi^2 v} \left( c_1 Z Z_{\mu\nu} F^{\mu\nu} h \right. - \left. c_2 Z (\partial_\mu h Z_\nu - \partial_\nu h Z_\mu) F^{\mu\nu} - \overline{\epsilon}_Z Z_{\mu\nu} \overline{F}^{\mu\nu} h \right),$$

where $\epsilon$ is the positron electric charge, $g$ is the $SU(2)_L$ coupling constant and $v = (\sqrt{2} G_F) ^{-1/2} \approx 246$ GeV is the vacuum expectation value of the Higgs field. Here $F_{\mu\nu}$ and $Z_{\mu\nu}$ are the standard field strengths for the electromagnetic and $Z$ field and $\overline{F}_{\mu\nu} = 2 F_{\mu\nu} - \epsilon_{\mu\nuab} F^{ab}/2$, with convention $\epsilon_{0123} = +1$. Dimensionless operators $c_1$, $c_2 Z$, $c_1 Z$, $c_1$, and $\overline{\epsilon}_Z$ are effective coupling constants $\overline{\epsilon}_Z$. As these coupling constants are, in general, complex-valued, the operators $\overline{\epsilon}_Z$ and $\overline{\epsilon}_Z$ are non-Hermitian, while being local and Lorentz invariant.

It is convenient to write the couplings $c_1$ and $c_1 Z$ as the sums of terms in the SM and new physics (NP) beyond the SM: $c_1 = c_1^{\text{SM}} + c_1^{\text{NP}}$, $c_1 Z = c_1 Z^{\text{SM}} + c_1 Z^{\text{NP}}$. In the SM, $\overline{\epsilon}_Z = c_2 Z = c_1 Z = 0$ and their nonzero values come from effects of the NP. The couplings $c_1^{\text{SM}}$ and $c_1 Z^{\text{SM}}$ have small imaginary parts which arise due to the intermediate on mass shell $\ell^+ \ell^- q \bar{q}$ states in the one-loop contributions [where $\ell = e, \mu, \tau$ denote leptons and $q = u, d, s, c, b$ denote quarks (excluding $t$ quark)]. In the one-loop order $c_1^{\text{SM}}$ and $c_1 Z^{\text{SM}}$ are given by \cite{15,20,21,22}

$$c_1^{\text{SM}} = A_1^f (\tau_f) + \sum_f N_f Q_f^2 A_1^{1/2} (\tau_f) \approx -6.60 + 0.08 i, \quad (4)$$

$$c_1 Z^{\text{SM}} = -A_2^f (\tau_f, \lambda_f) - \sum_f N_f Q_f g_f A_2^{1/2} (\tau_f, \lambda_f) \approx -5.540 + 0.005 i, \quad (5)$$
where $f = (\ell, q, t), N_f = 1(3)$ for leptons (quarks), $Q_f$ is the charge of the fermion $f$ in units of the electric charge of the positron. Here also $g_f = (2t_{\ell L} f - 4 Q_f \sin^2 \theta_W)/ \cos \theta_W$, where $t_{\ell L} f$ is the projection of the weak isospin of the $f$ fermion, and $\theta_W$ is the weak angle. The one-loop functions $A_f^1$, $A_f^2$, $A_f^3$, $A_f^3$ are defined in the Appendix A. These functions depend on arguments $\tau_W = 4 m_W^2/m^2, \lambda_W = 4 m_W^2/m^2, \tau_f = 4 m_f^2/m^2, \lambda_f = 4 m_f^2/m^2$, with $m_h$ being the mass of the Higgs boson, $m_W/m_Z$ being the mass of the $W$ ($Z$) boson, and $m_f$ being the mass of the $f$-th fermion. Numerical values in (13), (14) are obtained for $m_h = 126$ GeV using the SM parameters from [24], and the quark masses are chosen according to [24].

The terms $c_\gamma, c_{1Z}, c_{2Z}$ above correspond to a CP-even scalar $h$, while the terms $\tilde{c}_\gamma, \tilde{c}_Z$ indicate a CP-odd pseudoscalar $h$. The presence of both sets of terms means that $h$ is not a CP eigenstate. Interference of these terms lead to CP violating effects which reveal in polarization states of the photon. Generally, the couplings $c^\gamma, c_{1Z}, c_{2Z}, \tilde{c}_\gamma, \tilde{c}_Z$ may be complex.

The SM can be considered as effective low-energy theory of an underlying unknown theory at a scale $\Lambda$ (characteristic scale of the NP) which is much higher than the electroweak scale $v$. In effective field-theory language [24, 27, 31], the couplings $c^\gamma, c_{1Z}, c_{2Z}, \tilde{c}_\gamma, \tilde{c}_Z$ can be obtained from gauge invariant dimension-6 operators such as

$$
\mathcal{O}_B = \frac{g}{\Lambda^2} (D_{\mu} H)^\dagger (D_{\nu} H) B^{\mu \nu},
$$

$$
\mathcal{O}_W = \frac{g}{\Lambda^2} (D_{\mu} H)^\dagger \tau_k (D_{\nu} H) W^{\mu \nu}_k,
$$

$$
\mathcal{O}_{BB} = \frac{g^2}{2 \Lambda^2} H^\dagger H B_{\mu \nu} B^{\mu \nu},
$$

$$
\mathcal{O}_{BB} = \frac{g^2}{2 \Lambda^2} H^\dagger H B_{\mu \nu} B^{\mu \nu},
$$

$$
\mathcal{O}_{WW} = \frac{g^2}{2 \Lambda^2} H^\dagger H W_{\mu \nu} W^{\mu \nu}_k,
$$

$$
\mathcal{O}_{WW} = \frac{g^2}{2 \Lambda^2} H^\dagger H W_{\mu \nu} W^{\mu \nu}_k,
$$

$$
\mathcal{O}_{WB} = \frac{g^2}{2 \Lambda^2} H^\dagger H \tau_k W_{\mu \nu} B^{\mu \nu}_k,
$$

$$
\mathcal{O}_{WB} = \frac{g^2}{2 \Lambda^2} H^\dagger H \tau_k W^{\mu \nu}_k \bar{B}^{\mu \nu}_k.
$$

Here, $g'$ is the weak hypercharge gauge coupling, $B_{\mu \nu}$ is the field strength tensor for the hypercharge gauge group, $W^{\mu \nu}_k$ is the field strength tensor for the weak $SU(2)$ gauge group ($k = 1, 2, 3$), $H$ represents the Higgs doublet, and $\tau_k$ are the Pauli matrices for weak isospin. The operators $\mathcal{O}_i$ are CP even, and $\mathcal{O}_j$ are CP odd. The dual field-strength tensors are defined by $\tilde{X}_{\mu \nu} = (1/2) \varepsilon_{\mu \nu \alpha \beta} X^{\alpha \beta}$, for $X = B, W_k$. The corresponding effective Hamiltonian is

$$
H^{(6)}_{\text{eff}} = - \mathcal{L}^{(6)}_{\text{eff}} = \sum_i c_i \mathcal{O}_i + \sum_j \tilde{c}_j \mathcal{O}_j,
$$

where $i = (B, W, BB, WW, WB)$ and $j = (BB, WW, WB)$. The $h \gamma \gamma$ and $h \gamma Z$ couplings follow from the effective Lagrangian (7) by making the replacement $H \to (0, (v + h)/\sqrt{2})^T$ in the unitary gauge,
\[ c_{1Z}^{\text{NP}} = - \sum_f N_f s_f Q_f g_f A_{ij}^{\lambda_f}(\tau_f, \lambda_f) \]
\[ \approx 0.3253 s_b - (8.2 s_b + 1.2 s_c + 0.2 s_\tau) \times 10^{-3} \]
\[ + i(4.8 s_b + 0.5 s_c + 0.1 s_\tau) \times 10^{-3}, \]

(16)

\[ \tilde{c}_Z = - \sum_f N_f p_f q_f g_f I_2(\tau_f, \lambda_f) \]
\[ \approx -0.4939 p_b + (9.6 p_b + 1.3 p_c + 0.3 p_\tau) \times 10^{-3} \]
\[ - i(4.9 p_b + 0.5 p_c + 0.1 p_\tau) \times 10^{-3}, \]

(17)

where one-loop functions \( f(\tau_f), I_2(\tau_f, \lambda_f) \) are specified in the Appendix [A] and their arguments \( \tau_f, \lambda_f \) are defined after Eq. [3].

In obtaining the numerical values in (13)–(17) we have taken into account dominant contributions from the charm, bottom, top quarks and \( \tau \) lepton, in particular, the charm, bottom quarks and \( \tau \) lepton give rise to the imaginary parts of the couplings in (13)–(17).

In terms of the parameters \( s_f \) and \( p_f \) the width of the decay \( h \to ff \) is written as
\[ \Gamma(h \to ff) = \frac{N_f G_F m_f^2 m_h \beta_f \left((1 + s_f)^2 \beta_f^2 + p_f^2\right)}{4 \sqrt{2} \pi}, \]

(18)

where \( \beta_f = \sqrt{1 - 4m_f^2/m_h^2} \) is velocity of fermion \( f = (\ell, q) \) in the rest frame of \( h \). With a good accuracy one can put \( \beta_f = 1 \). Note that if one chooses \((1 + s_f)^2 + p_f^2 = 1\), then the width in Eq. (18) coincides with the decay width of the SM Higgs boson.

III. AMPLITUDES AND ANGULAR DISTRIBUTIONS

Let us consider the decay of the zero-spin Higgs boson into a pair of photons
\[ h(p) \to \gamma(k_1, \epsilon_1) \gamma(k_2, \epsilon_2), \]

(19)

where \( p \) is the four-momentum of \( h \) boson, \( k_1, k_2 \) are the four-momenta of photons and \( \epsilon_1, \epsilon_2 \) are the corresponding polarization four-vectors. In the rest frame of \( h \), the amplitude of this decay can be written in the form
\[ A(h \to 2\gamma) = \frac{e^2 m_h^2}{16 \pi^2 v} \left( c_\gamma (\epsilon_1^* \epsilon_2^2) + \tilde{c}_\gamma (\tilde{k} | \epsilon_1^* \times \epsilon_2^2) \right), \]

(20)

where \( m_h \) is the mass of \( h \) boson. The polarization vectors are chosen in the form \( \epsilon_1 = (0, \epsilon_1^2), \epsilon_2 = (0, \epsilon_2^2) \), where \( \tilde{c}_1 \tilde{k} = \tilde{\epsilon}_2 \tilde{k} = 0, \tilde{k} \) is the three-momentum of one of the photons and \( \tilde{k} = \tilde{k}/|\tilde{k}| \).

The helicity amplitudes for decay (19) are equal to
\[ H_\pm = -\frac{e^2 m_h^2}{16 \pi^2 v} (c_\gamma \pm i \tilde{c}_\gamma). \]

(21)

The decay width of \( h \to 2\gamma \) is
\[ \Gamma(h \to 2\gamma) = \frac{1}{32 \pi m_h} \left(|H_+|^2 + |H_-|^2\right). \]

(22)

The polarization states of a single photon are usually described through the density matrix \( \rho^{(\gamma)} \). For the process (19), one can write the two-photon density matrix following Ref. [34] as follows:
\[ \rho^{(\gamma \gamma)} = \frac{1}{4} \left( 1 - \sigma_3 \otimes \sigma_3 + \xi_1 (\sigma_1 \otimes \sigma_2 - \sigma_2 \otimes \sigma_1) + \xi_2 (\sigma_3 \otimes 1 - 1 \otimes \sigma_3) - \xi_3 (\sigma_1 \otimes \sigma_1 + \sigma_2 \otimes \sigma_2) \right), \]

(23)

where \( \sigma = (\sigma_1, \sigma_2, \sigma_3) \) are the Pauli matrices, \( 1 \) is 2 x 2 unit matrix, and \( \otimes \) means the direct product of two matrices. The reference frame is chosen with the \( Z \) axis along \( \tilde{k} \), and matrices on the left (right) from symbol \( \otimes \) refer to the photon with momentum \( \tilde{k} \) (\(-\tilde{k}\)).

In (23) the following parameters are introduced
\[ \xi_1 = \frac{2 \text{Im} (H_+ H_+^*)}{|H_+|^2 + |H_-|^2} = \frac{2 \text{Re}(c_\gamma \tilde{c}_\gamma^*)}{|c_\gamma|^2 + |\tilde{c}_\gamma|^2}, \]
\[ \xi_2 = \frac{|H_+|^2 - |H_-|^2}{|H_+|^2 + |H_-|^2} = \frac{2 \text{Im}(c_\gamma \tilde{c}_\gamma^*)}{|c_\gamma|^2 + |\tilde{c}_\gamma|^2}, \]
\[ \xi_3 = -\frac{2 \text{Re}(H_+ H_-^*)}{|H_+|^2 + |H_-|^2} = \frac{|\tilde{c}_\gamma|^2 - |c_\gamma|^2}{|c_\gamma|^2 + |\tilde{c}_\gamma|^2}. \]

The Stokes parameter \( \xi_2 \) defines degree of circular polarization of the photon with momentum \( \tilde{k} \), it has the meaning of average photon helicity. Parameters \( \xi_1, \xi_3 \) define correlation of linear polarizations of two photons (in particular, for \( \xi_1 = 0, \xi_3 = -1 \) the linear polarizations are parallel, while for \( \xi_1 = 0, \xi_3 = 1 \) they are orthogonal).

Next we come to the decay of \( h \) to \( \gamma \) and \( Z \) boson
\[ h(p) \to \gamma(k_1, \epsilon_1) Z(k_2, \epsilon_2), \]

(25)

where \( k_1, (k_2) \) is the four-momentum of photon (\( Z \) boson), \( \epsilon_1, (\epsilon_2) \) is polarization vector of the photon (\( Z \) boson).

The helicity amplitudes for the decay (26) are
\[ H_\pm = -\frac{eg m_h^2}{16 \pi^2 v} \left( 1 - \frac{m_\gamma^2}{m_h^2} \right) (c_{1Z} + c_{2Z} \pm i \tilde{c}_Z), \]

(26)

with the decay width
\[ \Gamma(h \to \gamma Z) = \frac{1}{16 \pi m_h} \left( 1 - \frac{m_\gamma^2}{m_h^2} \right) \left(|H_+|^2 + |H_-|^2\right), \]

(27)

where \( m_\gamma \) is the \( Z \) boson mass.

From definitions (24) we find the polarization parameters
\[ \xi_1 = -\frac{2 \text{Im}(A_{11} A_{11}^*)}{|A_{11}|^2 + |A_{12}|^2}, \]
\[ \xi_2 = \frac{2 \text{Re}(A_{11} A_{11}^*)}{|A_{11}|^2 + |A_{12}|^2}, \]
\[ \xi_3 = \frac{|A_{11}|^2 - |A_{12}|^2}{|A_{11}|^2 + |A_{12}|^2}. \]
where $H_\pm$ from Eq. (26) for further convenience are replaced by the amplitudes $A_f = (H_+ + H_-)/\sqrt{2}$ and $A_{\perp} = (H_+ - H_-)/\sqrt{2}$ corresponding to linearly polarized final states.

Numerical values of parameters $\xi_1$, $\xi_2$, $\xi_3$ will be discussed in Sec. IV.

In the decay $^{(23)}$, due to the zero-spin nature of the Higgs boson, the photon and $Z$ boson have equal helicities. This allows for measurement of the photon circular polarization through the decay $h \rightarrow \gamma Z \rightarrow \gamma f \bar{f}$ $^{(32)}$. Indeed, we derive the following angular distribution of the process in the polar angle $\theta$ between the momentum of the fermion $f$ in the $Z$ boson rest frame and the direction of the $Z$ boson motion in the $h$ boson rest frame,

$$\frac{1}{\Gamma} \frac{d\Gamma(h \rightarrow \gamma Z \rightarrow \gamma f \bar{f})}{d\cos\theta} = \frac{3}{8} \left(1 + \cos^2 \theta - 2 A^{(f)} \xi_2 \cos \theta\right), \quad (29)$$

where

$$A^{(f)} = \frac{2 g^f_{\nu} g^f_A}{(g^f_{\nu})^2 + (g^f_A)^2}. \quad (30)$$

The vector $g^f_{\nu}$ and axial-vector $g^f_A$ constants are

$$g^f_{\nu} \equiv t_{3L,f} - 2 Q_f \sin^2 \theta_W, \quad g^f_A \equiv t_{3L,f}. \quad (31)$$

Measurement of the forward-backward asymmetry $A_{FB}$ relative to the direction of $Z$ boson motion in the $h$ boson rest frame for the $f$ fermions produced in decay $h \rightarrow \gamma Z$

$$A_{FB} = \frac{F - B}{F + B}, \quad (32)$$

where

$$F \equiv \int_0^1 \frac{1}{\Gamma} \frac{d\Gamma}{d\cos\theta} d\cos\theta, \quad B \equiv \int_{-1}^0 \frac{1}{\Gamma} \frac{d\Gamma}{d\cos\theta} d\cos\theta,$$

which is

$$A_{FB} = -\frac{3}{4} A^{(f)} \xi_2, \quad (33)$$

allows one to find $\xi_2$.

Note that $A^{(\nu)}$ for the decay

$$h \rightarrow \gamma Z \rightarrow \gamma \mu^- \mu^+ \quad (34)$$

is $0.142 \pm 0.015 \quad (35)$, therefore in view of the condition $|\xi_2| \leq 1$, the absolute value of the asymmetry for this decay is not larger than 0.11. At the same time for the decay channel

$$h \rightarrow \gamma Z \rightarrow \gamma b \bar{b} \quad (35)$$

($A^{(b)} = 0.923 \pm 0.020 \quad (35)$), the absolute value of $A_{FB}$ can be much larger, namely, as large as 0.69.

Consider now feasibility to measure the distribution $^{(29)}$ at the LHC after its upgrade to higher luminosity and energy $\sqrt{s} = 14$ TeV. Taking into account various mechanisms of Higgs boson production in $pp$ collisions, the inclusive cross section is $\sigma = 57.0163 \text{ pb} \quad (20)$. Then the cross section for the process $pp \rightarrow h X \rightarrow \gamma \ell^+ \ell^- X$ in the SM is

$$\sigma \times \text{BR}(h \rightarrow \gamma Z) \text{BR}(Z \rightarrow \ell^+ \ell^-) = 6.24 \text{ fb}, \quad (36)$$

with $\ell = e, \mu$ and the branching fractions are taken from Refs. $^{23} \quad (30)$. In order to observe the forward-backward asymmetry $A_{FB}$ for maximal value $|\xi_2| = 1$ at a $3 \sigma$ level, the number of events should be bigger than 734. This number of events can be obtained, with ideal detector, with integrated luminosity about 120 fb$^{-1}$.

Let us discuss a possibility to determine the polarization parameters $\xi_1$ and $\xi_3$. For this one can study the process

$$h \rightarrow \gamma^\ast Z \rightarrow \ell^+ \ell^- Z \quad (37)$$

with the decay $Z \rightarrow \bar{f} f$ on mass shell. For the process $^{(37)}$ we obtain the distribution over the dilepton invariant mass squared $q^2$ and azimuthal angle $\phi$ between the decay planes of $\gamma^\ast \rightarrow \ell^+ \ell^-$ and $Z \rightarrow \bar{f} f$ in the $h$ rest frame:

$$\frac{d\Gamma(h \rightarrow \ell^+ \ell^- Z)}{dq^2 d\phi} = \frac{1}{2\pi} \left(1 - \frac{1}{4} (1 - F_L(q^2)) \times \left(\xi_3(q^2) \cos 2\phi + \xi_1(q^2) \sin 2\phi\right)\right). \quad (38)$$

Here

$$F_L(q^2) \equiv \frac{|A_0(q^2)|^2}{|A_0(q^2)|^2 + |A^{(\nu)}(q^2)|^2 + |A_{\perp}(q^2)|^2} \quad (39)$$

is the fraction of longitudinal polarization of virtual photon, and the amplitudes are defined as

$$A_0(q^2) = \frac{e g}{16 \pi^2 v} \sqrt{\frac{q^2}{m_H^2}} \left(2 c_{1Z} m_Z^2 + c_{2Z} (m_H^2 - q^2 + m_Z^2)\right), \quad (40)$$

$$A^{(\nu)}(q^2) = -\frac{e g}{8 \sqrt{2} \pi^2 v} \left(c_{1Z} (m_H^2 - q^2 - m_Z^2) + c_{2Z} (m_H^2 + q^2 - m_Z^2)\right), \quad (41)$$

$$A_{\perp}(q^2) = -i \frac{e g}{8 \sqrt{2} \pi^2 v} \bar{c}_Z \sqrt{\lambda(m_H^2, q^2, m_Z^2)}, \quad (42)$$

with $\lambda(a,b,c) \equiv a^2 + b^2 + c^2 - 2(ab + ac + bc)$ and the distribution over the invariant mass squared reads

$$\frac{d\Gamma}{dq^2} = \frac{\alpha_{em} \sqrt{\lambda(m_H^2, q^2, m_Z^2)}}{48 \pi^2 m_H^2 q^2} \left(|A_0(q^2)|^2 + |A^{(\nu)}(q^2)|^2 + |A_{\perp}(q^2)|^2\right), \quad (43)$$
where \( \alpha_{em} = e^2/(4\pi) \) is the electromagnetic fine-structure constant. The \( q^2 \)-dependent quantities \( \xi_1(q^2) \) and \( \xi_2(q^2) \) can be obtained from Eqs. (23) in which the amplitudes \( A_\| (A_\perp) \) are substituted by the \( q^2 \)-dependent amplitudes \( A_\| (q^2) \) (\( A_\perp (q^2) \)). In derivation of (25) we assumed that leptons are massless.

In expressions (44)–(46) we did not take into account additional two-fermion current operators of dimension 6 in the effective Hamiltonian (7) and the process \( h \to Z^+ Z \to \ell^+ \ell^- Z \). Both these mechanisms contribute to tree level to the decay \( h \to \ell^+ \ell^- Z \).

From (48) one can approximately find \( \xi_1 \) and \( \xi_3 \) in the decay \( h \to \gamma Z \). Neglecting the amplitude (40) for longitudinally polarized photon \( A_\| (q^2) \approx q^2 \), and \( q^2 \)-dependence of the transverse amplitudes, i.e. substituting \( A_\| (q^2) \approx A_\| (0) \) and \( A_\perp (q^2) \approx A_\perp (0) \), we obtain the distribution over the azimuthal angle

\[
\frac{d\Gamma(h \to \ell^+ \ell^- Z)}{d\phi} \approx \left( \frac{\alpha_{em}}{3\pi} \log \frac{q^2_{\text{max}}}{q^2_{\text{min}}} \right) \Gamma(h \to \gamma Z) \times \frac{1}{2\pi} \left( 1 - \frac{1}{4} (\xi_1 \cos 2\phi + \xi_1 \sin 2\phi) \right). \tag{44}
\]

The lower integration limit \( q^2_{\text{min}} \) is determined by possibilities of detectors, in particular, to provide sufficient \( \phi \) resolution to separate \( \sin 2\phi \) and \( \cos 2\phi \) terms in the distribution (44). In this connection we should mention recent measurements of the \( B^0 \to K^{*0} e^+ e^- \) branching fraction \( \mathcal{B}(K^{*0} e^+ e^-) \), in which the LHCb detector allowed selection of the lower value of dilepton invariant mass equal to 30 MeV.

Theoretical accuracy of Eq. (44) improves with the decreasing value of \( q^2_{\text{max}} \), since contribution of the competing mechanism \( h \to Z^* Z \to \ell^+ \ell^- Z \) diminishes for \( q^2_{\text{max}} \ll m_Z^2 \). Consider for example production of the \( e^+ e^- \) pair in the process \( h \to e^+ e^- Z \) with dilepton invariant mass from 30 MeV to 1000 MeV. Our calculation including both \( h \to \gamma Z \to e^+ e^- Z \) and \( h \to Z^* Z \to e^+ e^- Z \) amplitudes shows that theoretical error in \( \xi_1, \xi_3 \), which arises when neglecting the \( h \to Z^* Z \to e^+ e^- Z \) mechanism, amounts to 20% in the SM (in which \( \xi_3^M = 0, \xi_3^P = -1 \)), and 10% in the effective Hamiltonian approach \( \mathcal{H} \) [the choice of coefficients \( \xi_1, \xi_3 \)] is discussed in Sec. IV.

Of course, the process \( h \to \gamma^* Z \to e^+ e^- Z \) is rare. Let us make an estimate of its observability at the LHC energy \( \sqrt{s} = 14 \text{ TeV} \). Using (44) and choosing the Higgs production inclusive cross section \( \sigma = 57.0163 \text{ pb} \) we calculate the SM cross section for the \( p p \to h X \to \gamma^* Z X \to e^+ e^- Z X \) in the interval of dilepton invariant mass from 30 MeV to 1000 MeV,

\[
\sigma \times \frac{\Gamma(h \to e^+ e^- Z)|_{30 < m_{ee} < 1000 \text{ MeV}}}{\Gamma(h \to \text{all})} = 0.5 \text{ fb}. \tag{45}
\]

When detecting \( Z \) boson via \( Z \to e^+ e^- \) and \( Z \to \mu^+ \mu^- \) channels the cross section (45) is reduced by factor 0.067, and for the integrated luminosity of 100 fb\(^{-1}\) we can expect about 3 events. This number is too small and a higher integrated luminosity will be needed to observe the decay \( h \to \gamma^* Z \to \ell^+ \ell^- Z \) and analyze its angular distribution.

### IV. RESULTS OF CALCULATION AND DISCUSSION

First we note that in the SM the polarization parameters are \( \xi_1^S = \xi_2^S = 0 \) and \( \xi_3^S = -1 \). Any deviations of the measured values of \( \xi_i \) from \( \xi_i^S \) \( (i = 1, 2, 3) \) will indicate presence of effects beyond the SM.

In order to estimate magnitude of effects of NP, we consider (i) the approach in which NP is expressed through dimension-6 operators described by effective Hamiltonian (7), and (ii) the model (13) with the scalar and pseudoscalar couplings of fermions to the Higgs boson.

In the approach (7) we take for definiteness \( c_W = c_W = c_W = 1 \), \( c_{WW} = c_{BB} = c_{BB} = c_{WW} = 1 \), choosing the scale \( \Lambda = 4\pi v \approx 3.1 \) TeV we obtain for the \( h \to \gamma \gamma \) decay

\[
\xi_1 = -0.259, \quad \xi_2 = 0.003, \quad \xi_3 = -0.966,
\]

\[
\mu_{\gamma \gamma} \equiv \frac{\Gamma(h \to \gamma \gamma)}{\Gamma_{\text{SM}}(h \to \gamma \gamma)} = 1.35, \tag{46}
\]

and for \( h \to \gamma Z \) decay

\[
\xi_1 = -0.107, \quad \xi_2 = 0.0001, \quad \xi_3 = -0.994,
\]

\[
\mu_{\gamma Z} \equiv \frac{\Gamma(h \to \gamma Z)}{\Gamma_{\text{SM}}(h \to \gamma Z)} = 1.12. \tag{47}
\]

For another scale \( \Lambda = 2 \text{ TeV} \), for the \( h \to \gamma \gamma \) decay, we obtain

\[
\xi_1 = -0.497, \quad \xi_2 = 0.004, \quad \xi_3 = -0.868,
\]

\[
\mu_{\gamma \gamma} = 1.99, \tag{48}
\]

and for \( h \to \gamma Z \) decay

\[
\xi_1 = -0.236, \quad \xi_2 = 0.0002, \quad \xi_3 = -0.972,
\]

\[
\mu_{\gamma Z} = 1.31. \tag{49}
\]

For the ratio \( \mu_{\gamma \gamma} \) our calculation with the scale \( \Lambda = 4\pi v \) better agrees with the ATLAS data (10) for \( h \to \gamma \gamma \) than calculation with \( \Lambda = 2 \text{ TeV} \).

In the model with scalar and pseudoscalar couplings of fermions to the Higgs boson (13) we choose the parameters

\[
p_t = p_b = p_c = p_r = \pm \sqrt{2},
\]

\[
s_t = s_b = s_c = s_r = 1/\sqrt{2} - 1 \tag{50}
\]

satisfying normalization \((1 + s_f)^2 + p_f^2 = 1\) discussed in Sec. III.

As a result, for the decay \( h \to \gamma \gamma \) we find

\[
\xi_1 = \mp 0.528, \quad \xi_2 = \mp 0.010, \quad \xi_3 = -0.849,
\]

\[
\mu_{\gamma \gamma} = 1.26 \tag{51}
\]
and for decay $h \to \gamma Z$

$$\xi_1 = \pm 0.121, \quad \xi_2 = \mp 0.001, \quad \xi_3 = -0.993,$$

$$\mu_{\gamma Z} = 1.04. \quad (52)$$

In addition, the $h \to f \bar{f}$ decay width calculated with $s_f, p_f$ in $[50]$ coincides with the SM decay width and agrees with the CMS data $[11]$ for $h \to \tau^+ \tau^-$ and $h \to b \bar{b}$ decays,

$$\mu_{\tau\tau} = \frac{\Gamma(h \to \tau^+ \tau^-)}{\Gamma_{SM}(h \to \tau^+ \tau^-)} = 1.10 \pm 0.41,$$

$$\mu_{bb} = \frac{\Gamma(h \to b \bar{b})}{\Gamma_{SM}(h \to b \bar{b})} = 1.15 \pm 0.62. \quad (53)$$

At the same time the channel $h \to c \bar{c}$ is not measured yet. Thus the $h \to c \bar{c}$ width, in general, may differ from the SM prediction, and consequently the constraint $(1 + s_c)^2 + p_c^2 = 1$ for the charm quark may not hold. We can make an assumption that $\Gamma(h \to c \bar{c}) \leq \Gamma(h \to b \bar{b})$. Combining this inequality with Eqs. (18) and (53) we find

$$(1 + s_c)^2 + p_c^2 \leq \mu_{bb} \times \frac{\Gamma_{SM}(h \to b \bar{b})}{\Gamma_{SM}(h \to c \bar{c})}. \quad (54)$$

Taking the central values of $\mu_{bb}$ and the widths from $[36]$ (Table 1 therein) we obtain the following constraint for the $c \bar{c}$ couplings: $(1 + s_c)^2 + p_c^2 \leq 22.8$.

To estimate maximal values of polarization parameter $\xi_2$ in the channel $h \to \gamma Z$ let us take $s_c, p_c$ satisfying $(1 + s_c)^2 + p_c^2 = 22.8$, although the latter equality does not fix $s_c, p_c$ uniquely. In addition, put $s_f = p_f = 0$ for $f \neq c$. Then calculation using $[16]$ and $[17]$ gives values of $\xi_2$ which do not exceed $8.6 \times 10^{-4}$. It is seen that even for such a radical modification of the Higgs couplings to the charm quarks, the parameter $\xi_2$ remains very small. Thus the existing data on the Higgs boson decay to the $\tau^+ \tau^-$ and $b \bar{b}$ pairs and a reasonable assumption on the upper bound of the decay width to the charm quarks lead to conclusion that the rescattering effects on the one-loop level result in values of $\xi_2$ in the $h \to \gamma Z$ decay about $10^{-3}$ or smaller.

It would be of interest to check in the experimental analysis of the distribution $[29]$ whether the parameter $\xi_2$ is very small indeed. If the analysis yielded sizable values of $\xi_2$, this would mean the presence of additional sources of non-Hermiticity of effective Lagrangian. The latter may arise, for example, due to the breaking of Hermiticity in an underlying (fundamental) theory at very small distances. Note, that similar aspects have been discussed in $[38]$ for the process $\gamma \gamma \to h$, where the authors calculated various asymmetries as functions of complex coefficients $c_{\gamma}, c_0$ in Eq. (2). Since the requirement of Hermiticity is one of the conditions in the proof of the $CPT$ theorem $[16]$, measurement of the photon circular polarization in the decay $h \to \gamma Z \to \gamma f \bar{f}$ through the forward-backward asymmetry $A_{FB}$ can be useful for testing $CPT$ symmetry.

The parameters $\xi_1$ and $\xi_3$ carry information on the $CP$ properties of the Higgs boson. Besides, $\xi_1$ is $CPT$-odd and $T$-odd observable and, in the absence of final-state interaction between the leptons and fermions, a nonzero value of $\xi_1$ will point to the violation of $T$ invariance.

V. CONCLUSIONS

In this paper polarization properties of the $\gamma \gamma$ and $\gamma Z$ states in the decays $h \to \gamma \gamma$ and $h \to \gamma Z$ of recently discovered scalar boson have been considered. We have chosen effective Lagrangian, describing $h \gamma \gamma$ and $h \gamma Z$ interactions with $CPT$-even and $CP$-odd parts. This allowed for calculation of polarization parameters $\xi_1$, $\xi_2$, $\xi_3$. In the SM these parameters take on values $\xi_1^{SM} = \xi_2^{SM} = 0$, $\xi_3^{SM} = -1$ and deviations of the measured values of $\xi_i$ from $\xi_i^{SM} (i = 1, 2, 3)$ will point to effects of NP.

The parameter $\xi_2$, which defines correlation of linear polarizations of $\gamma$ and $Z$, can be extracted from the azimuthal angle distribution in the process $h \to \gamma^\pm Z \to \ell^+ \ell^- Z$ with decay $Z \to f \bar{f}$ on the mass shell.

In numerical estimates of these parameters we included the one-loop contribution from the SM, and models beyond the SM. Namely, we applied the approach $[24, 27–31]$ in which NP is described by dimension-6 operators in the fields of the SM, and model with scalar and pseudoscalar couplings of fermions to the Higgs boson on the one-loop level.

The value of photon circular polarization turns out to be very small, of the order $10^{-3}$. In general, nonzero value of $\xi_2$ arises due to presence of the $CP$-even and $CP$-odd parts in effective Lagrangian $L^{h\gamma\gamma}_{\text{eff}}$ and absorptive parts of one-loop diagrams, or rescattering effects of the type $h \to a a \to \gamma Z$, where $a$ are charged particles with masses $m_a \leq m_h/2$. Only leptons and quarks $u, d, s, c, b$ satisfy this condition and hence contribute to absorptive parts of one-loop diagrams. Contributions from leptons $e, \mu$ and light quarks $u, d, s$ are negligibly small. The couplings of $h$ to the $\tau$ lepton and bottom quark are constrained by recent CMS data on the $h \to \tau^+ \tau^-$ and $h \to b \bar{b}$ decays, and couplings to the charm quark are constrained from an assumption on the upper bound of the $h \to c \bar{c}$ decay width.

Apart from rescattering effects, in framework of $CPT$ symmetric models, there are no sources of non-Hermiticity of $L^{h\gamma\gamma}_{\text{eff}}$ which could contribute to parameter $\xi_2$. If there is a violation of $CPT$ symmetry in an underlying theory at small distances, then this may give rise to additional non-Hermiticity effects in $L^{h\gamma\gamma}_{\text{eff}}$ which will change the value of $\xi_2$. Therefore measurement of this parameter in the $h \to \gamma Z \to \gamma f \bar{f}$ process would allow one to test the prediction of the SM, and to search for deviations from the SM, and even possible effects of
CPT violation in an underlying theory.

Nonzero values of parameter $\xi_1$ point to violation of CP symmetry in the $h \to \gamma \gamma$ and $h \to \gamma Z$ decays. In the chosen models of NP, for the $h \to \gamma Z$ decay, $\xi_1$ appears to be 0.1-0.2. Its experimental determination can put constraints on models describing physics beyond the SM.

We also estimated in the SM a feasibility of measurement of the discussed processes in the $pp$ collisions at the LHC, after its upgrade to energy $\sqrt{s} = 14$ TeV and higher luminosity. The cross section for the process $pp \to h X \to \gamma Z X \to \ell^+\ell^-\ell^+\ell^- X$ ($\ell = e, \mu$) turns out to be 6.24 fb. With integrated luminosity about 120 fb$^{-1}$ and ideal detector it may be possible to observe the forward-backward asymmetry $A_{FB}$ for $|\xi_2| = 1$ at a 3$\sigma$ level.

Here we should mention papers [39, 40], where possibilities of studying at the LHC the $h \to \gamma \ell^+\ell^-$ decay via $\gamma Z$ channel are considered. Although observation of the Higgs is difficult in view of the background which is more rare process, and our estimate of its observability is less optimistic. One can expect about 3 events in the interval of $e^+e^-$ invariant mass from 30 MeV to 1000 MeV if $Z$ boson is detected through the $Z \to e^+e^-$, $\mu^+\mu^-$ channels. Clearly an integrated luminosity higher than 100 fb$^{-1}$ will be needed to study the $h \to \gamma^* Z \to e^+e^- Z$ process.

In conclusion, we hope that with increasing the integrated luminosity at the LHC investigation of angular distributions discussed in the present paper will become possible.

Appendix A: Definition of Loop Functions

The loop functions for the $W^\pm$ boson ($A_{1/2}^{(Z)}$) as well as the fermion $f$ ($A_{1/2}^{\gamma(Z)}$) are defined in Ref. [23]

$$A_{1/2}^{(Z)}(\tau) = - (2 + 3\tau + 3\tau(2 - \tau)f(\tau)) , \quad (A1)$$

$$A_{1/2}^{\gamma(Z)}(\tau) = 2\tau(1 + (1 - \tau)f(\tau)) , \quad (A2)$$

$$A_1^Z(\tau, \lambda) = \cos \theta_W \left( 4 \left( 3 - \tan^2 \theta_W \right) I_2(\tau, \lambda) + \left( \left( 1 + \frac{2}{\tau} \right) \tan^2 \theta_W - \left( 5 + \frac{2}{\tau} \right) \right) I_1(\tau, \lambda) \right) . \quad (A3)$$

$$A_1^{Z/2}(\tau, \lambda) = I_1(\tau, \lambda) - I_2(\tau, \lambda) . \quad (A4)$$

The functions $I_1, I_2$ are given by

$$I_1(\tau, \lambda) = \frac{\tau \lambda}{2(\tau - \lambda)} \left( 1 + \frac{\tau \lambda}{\tau - \lambda} (f(\tau) - f(\lambda)) \right) + \frac{2}{\tau - \lambda} (g(\tau) - g(\lambda)) , \quad (A5)$$

$$I_2(\tau, \lambda) = - \frac{\tau \lambda}{2(\tau - \lambda)} (f(\tau) - f(\lambda)) , \quad (A6)$$

where the functions $f(\tau)$ and $g(\tau)$ can be expressed as

$$f(\tau) = \begin{cases} \arcsin \frac{1}{\sqrt{\tau}} & \tau \geq 1 \\ \frac{1}{4} \left( \log \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} - i\pi \right)^2 & \tau < 1 \end{cases} . \quad (A7)$$

$$g(\tau) = \begin{cases} \sqrt{\tau - 1} \arcsin \frac{1}{\sqrt{\tau}} & \tau \geq 1 \\ \sqrt{1 - \tau} \arcsin \frac{1}{\sqrt{1 - \tau}} \left( \log \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} - i\pi \right) & \tau < 1 \end{cases} . \quad (A8)$$
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