Rare decay $Z\to\bar{\nu}\nu\gamma\gamma$ via quartic gauge boson couplings

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We present a detailed calculation of the rare decay $Z\to\bar{\nu}\nu\gamma\gamma$ via the quartic neutral gauge boson coupling $ZZ\gamma\gamma$ in the framework of the effective Lagrangian approach. The current experimental bound on this decay mode is then used to constrain the coefficients of this coupling. It is found that the bounds obtained in this way, of the order of 10^{-1}, are weaker than the ones obtained from the analysis of triple-boson production at the CERN e^-e^+ collider LEP-2.

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I. INTRODUCTION

One of the most sensitive probes of physics beyond the standard model (SM) is provided by neutral gauge boson self couplings [1]. In the SM, trilinear neutral gauge boson couplings (TNGBCs) $V_iV_jV_k$ and quartic neutral gauge boson couplings (QNGBCs) $V_iV_jV_kV_l$, with $V_i=Z,\gamma$, vanish at the tree level and their radiative corrections are known to be rather small, of the order of 10^{-8}-10^{-10} [2,3]. Deviations from the SM predictions for TNGBCs and QNGBCs might point to new interactions such as those arising from strongly interacting electroweak models [4], a fourth family of chiral fermions with SM assignments of quantum numbers [2], and the heavy fermions arising in the minimal supersymmetric standard model (MSSM) [3]. An interesting feature of QNGBCs involving at least one photon field stems from the fact that they are genuine in the sense that arise from effective operators that do not induce any trilinear gauge boson coupling. Therefore these genuine QNGBCs must be constrained from processes other than the ones used to constrain trilinear gauge boson couplings, such as boson-pair fusion or triple-boson production [5]. This is to be contrasted with the case of the quartic $WW\gamma\gamma$ coupling, which can arise from operators that also induce the trilinear $WW\gamma$ coupling, which means that any constraint on the latter can be immediately translated into a bound on the former. Furthermore, while the quartic neutral $ZZZZ$ coupling can be induced at the tree level, for instance by the exchange of a heavy scalar boson, QNGBCs involving at least one photon field can only arise at one-loop level or higher order in any renormalizable theory because of electromagnetic gauge invariance. In particular, QNGBCs have been constrained from the CERN e^+e^- collider LEP-2 data on triple gauge boson production [5]. Considerable work has also been devoted to the analysis of QNGBCs through different processes at the next linear e^+e^- collider as well as $\gamma\gamma$, $e\ \gamma$ and hadron colliders [4,6].

In this Brief Report we will present a detailed calculation of the $ZZ\gamma\gamma$ coupling contribution to the rare decay $Z\to\bar{\nu}\nu\gamma\gamma$. The current experimental limit on this decay mode [7,8] will then be used to constrain the coefficients of this QNGBC. To our knowledge, this calculation has never been presented in the literature. Our analysis will proceed in the same line as those presented in our previous works [9,10], where we obtained bounds on TNGBCs from the $Z\to\bar{\nu}\nu\gamma$ decay mode [9] and on neutrino-photon interactions from $Z\to\bar{\nu}\nu\gamma\gamma$ [10]. We will find that the bounds obtained in this way rely on very few assumptions, though they are weaker than the ones obtained from the analysis of $e^+e^-\to Z\gamma\gamma$ and $e^+e^-\to W^+W^-\gamma$ data at LEP-2 [5]. The organization of the paper is as follows. In Sec. II we present a short description of the effective Lagrangian for the $ZZ\gamma\gamma$ coupling. Section III is devoted to present the calculation of the rare decay $Z\to\bar{\nu}\nu\gamma\gamma$. Finally, our conclusions are presented in Sec. IV.

II. THE EFFECTIVE COUPLING $ZZ\gamma\gamma$

When parametrizing physics beyond the Fermi scale in a model-independent manner by means of the effective Lagrangian technique there are two alternatives, i.e. the underlying new physics can be assumed to be either of decoupled or nondecoupled nature [11]. In the decoupling scenario it is assumed that the Higgs mechanism is realized in nature, thereby requiring the existence of at least one (relatively) light Higgs boson. In this case the low-energy theory is renormalizable in the manner of Dyson, the decoupling theorem remains valid, and the effective Lagrangian is con-

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structed out of operators respecting the $SU_L(2)\times U_Y(1)$ symmetry linearly. In this scenario the virtual heavy physics effects cannot affect dramatically the low-energy processes: the impact of new physics might become important only in those processes that are absent or very suppressed within the SM [11]. On the other hand, another possibility (the nondecoupling scenario) arises if the Higgs boson is very heavy or does not exist at all. There follows that the low-energy theory is nonrenormalizable due to the absence of the Higgs boson. This class of new physics can be assumed to be responsible for the symmetry breaking of the electroweak sector. In this scenario the $SU_L(2)\times U_Y(1)$ symmetry is nonlinearly realized [11].

The lowest-dimension operators that induce the $ZZ\gamma\gamma$ coupling have dimension six (eight) in the nonlinear (linear) realization. This means that any new physics effects arising from this coupling are likely to become more evident in the nonlinear scenario since in the linear one they are suppressed by higher powers of the new physics scale Λ. Therefore, in this work we will concentrate on the nonlinear scenario. Furthermore, only those operators that respect the custodial $SU_C(2)$ and the discrete C and P symmetries will be considered. It turns out that the operators that violate the custodial symmetry are tightly constrained by the ρ parameter. In the nonlinear scenario there are fourteen dimension-six operators that induce the $ZZ\gamma\gamma$ coupling at tree level [12]. This was discussed to a large extent in Ref. [12] and we will not dwell on this issue here. We rather focus on the Lorentz structure induced for the $ZZ\gamma\gamma$ coupling. In the unitarity gauge ($U=1$), there are only two independent Lorentz structures for this coupling induced by dimension-six operators [12]:

$$\mathcal{L}_{ZZ\gamma\gamma}=-\frac{e^2}{16\Lambda^2c_W^2}a_0F_{\mu\nu}F^{\mu\nu}Z^\alpha Z_\alpha -\frac{e^2}{16\Lambda^2c_W^2}a_cF_{\mu\nu}F^{\mu\alpha}Z^\nu Z_\alpha, \tag{1}$$

where $F_{\mu\nu}=\partial_\mu A_\nu-\partial_\nu A_\mu$ and $Z_{\mu\nu}=\partial_\mu Z_\nu-\partial_\nu Z_\mu$. We have followed closely the notation introduced in [12]. Below we will proceed to compute the decay $Z\to\nu\bar{\nu}\gamma\gamma$ using the above effective interaction. The experimental limit on this decay will then be used to constrain the coefficients a_0/Λ^2 and a_c/Λ^2.

III. THE RARE DECAY $Z\to\bar{\nu}\nu\gamma\gamma$

We now turn to outline the calculation of the decay width for $Z\to\bar{\nu}\nu\gamma\gamma$, which is somewhat similar to the one presented in Ref. [13] for the rare decay $Z\to\bar{\nu}\nu AA$ in two Higgs doublet models, with A the CP-odd neutral scalar. In Ref. [10], the experimental bound on the rare $Z\to\bar{\nu}\nu\gamma\gamma$ decay was used to constrain the neutrino-photon interactions $\bar{\nu}\nu\gamma$ and $\bar{\nu}\nu\gamma\gamma$. Here we will make an analysis along the same lines but focus on the purely bosonic $ZZ\gamma\gamma$ coupling, which contributes to the $Z\to\nu\bar{\nu}\gamma\gamma$ decay through the Feynman diagram shown in Fig. 1.

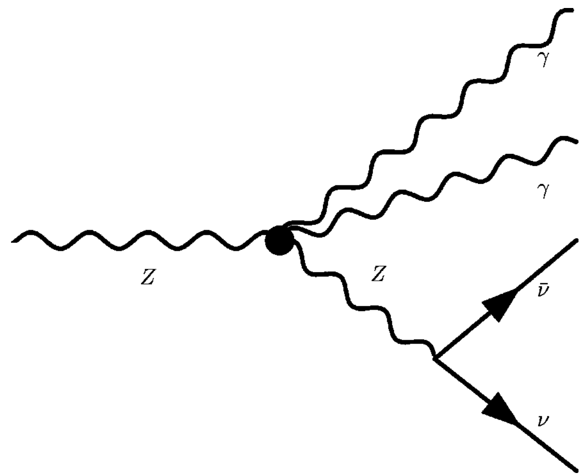

FIG. 1. Contribution from the effective $ZZ\gamma\gamma$ coupling to the rare $Z\to\bar{\nu}\nu\gamma\gamma$ decay.

The 4-vectors of the participating particles will be denoted as follows $Z(p)\to\bar{\nu}(p_1)\nu(p_2)\gamma(k_1)\gamma(k_2)$. The Feynman rule for the effective vertex $Z(q_\mu)Z(q'_\nu)\gamma(k_\alpha)\gamma(k'_\beta)$ is straightforwardly obtained from Eq. (1):

$$\frac{i\,e^2}{8\,\Lambda^2}\{4\,a_0g_{\mu\nu}(k\cdot k'g_{\alpha\beta}-k'_\alpha k_\beta) +a_c[k\cdot k'(g_{\mu\alpha}g_{\nu\beta}+g_{\mu\beta}g_{\nu\alpha})+g_{\alpha\beta}(k_\mu k'_\nu+k_\nu k'_\mu) -k_\beta(g_{\alpha\mu}k'_\nu+g_{\alpha\nu}k'_\mu)-k'_\alpha(g_{\beta\mu}k_\nu+g_{\beta\nu}k_\mu)]\} \tag{2}$$

where all the momenta are directed inward. The decay width can then be written as

$$\Gamma(Z\to\bar{\nu}\nu\gamma\gamma) =\frac{1}{(2\pi)^8 2^5 m_Z}\int |\bar{\mathcal{M}}|^2\delta^{(4)}\left(p-\sum_{i=1}^4 q_i\right)\prod_{i=1}^4\frac{d^3q_i}{2\,q_i^0}, \tag{3}$$

with $q_i=p_1$, p_2, k_1 and k_2 for $i=1$, 2, 3 and 4, respectively. By the usual method and after a lengthy calculation we can obtain the squared amplitude. It reads

$$|\bar{\mathcal{M}}|^2=\frac{2}{3}\left(\frac{g}{4c_W}\right)^2\frac{\Xi_{\mu\nu}p_1^\mu p_2^\nu}{[(p-k_1-k_2)^2-m_Z^2]^2}, \tag{4}$$

where we have introduced the definition

$$\Xi_{\mu\nu}=\left(\frac{\pi\,\alpha}{m_Z\Lambda^2c_W^2}\right)^2(|4\,a_0+a_c|^2(k_1\cdot k_2)^2A_{\mu\nu} +2\,|a_c|^2[m_Z^2(k_1\cdot k_2)+2k_1\cdot pk_2\cdot p]B_{\mu\nu}), \tag{5}$$

with $A_{\mu\nu}=m_Z^2g_{\mu\nu}+2\,p_\mu p_\nu$ and $B_{\mu\nu}=k_{1\mu}k_{2\nu}+k_{1\nu}k_{2\mu}$. In Eq. (4) a factor of 3 is included in the denominator as we are averaging over the Z boson polarizations; also, a factor of 2 has been introduced to account for two identical particles in the final state. The integration over p_1 and p_2 can be carried out straightforwardly with the aid of the following result [14]:

$$I^{\mu\nu}=\int \frac{d^3p_1}{p_1^0}\frac{d^3p_2}{p_2^0}\delta^{(4)}(Q-p_1-p_2)p_1^{\mu}p_2^{\nu}$$
$$=\frac{\pi}{6}(Q^2g^{\mu\nu}+2Q^{\mu}Q^{\nu}), \tag{6}$$

where $Q=p-p_1-p_2$. Once we are done with the integration over p_1 and p_2, there still remains to integrate over k_1 and k_2. To this end we will work in the center of mass frame of the Z boson. We find it useful to define the following variables $\xi=2\,p\cdot k_1/m_Z^2=2\,k_1^0/m_Z$, $\eta=2\,p\cdot k_2/m_Z^2=2\,k_2^0/m_Z$, and $\omega=(1-\cos\theta)/2$. The decay width can thus be expressed as

$$\Gamma(Z\to\bar{\nu}\nu\gamma\gamma)=\left(\frac{m_Z^2\alpha}{\Lambda^2c_W^2}\right)^2\frac{m_Z\alpha}{2(2^83\pi c_Ws_W)^2}$$
$$\times\int_{\Omega}h(\xi,\eta,\omega)d\xi d\eta d\omega, \tag{7}$$

where

$$h(\xi,\eta,\omega)=\frac{\xi^3\eta^3}{(\xi\eta\omega-\xi-\eta)^2}(|4a_0+a_c|^2f(\xi,\eta,\omega)$$
$$+|a_c|^2g(\xi,\eta,\omega)), \tag{8}$$

$$f(\xi,\eta,\omega)=\omega^2(12+\xi^2-\eta(12-\eta)-2\xi(6-\eta-4\omega\eta)), \tag{9}$$

and

$$g(\xi,\eta,\omega)=4(1+\omega)(1+\omega-2\omega\xi-2\eta\omega(1-\omega\xi)). \tag{10}$$

It can be shown that the integration region Ω is given by [14]

$$0\leqslant\omega\leqslant1$$

when

$$0\leqslant\xi\leqslant1-\eta, \tag{11a}$$

$$\frac{\eta+\xi-1}{\eta\xi}\leqslant\omega\leqslant1$$

when

$$1-\eta\leqslant\xi\leqslant1, \tag{11b}$$

together with $0\leqslant\eta\leqslant1$. After numerical integration we are left with

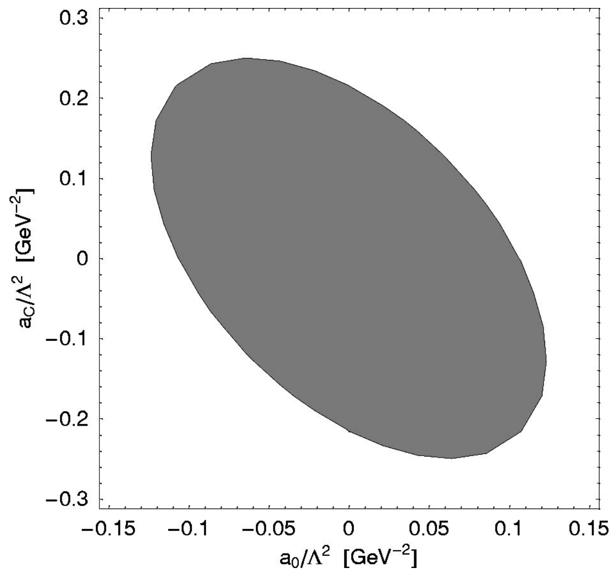

FIG. 2. Allowed area (gray region) in the a_0/Λ^2 vs a_c/Λ^2 plane from the experimental bound on the $Z\to\bar{\nu}\nu\gamma\gamma$ rare decay.

$$\Gamma(Z\to\bar{\nu}\nu\gamma\gamma)=\left(\frac{1\ \text{GeV}}{\Lambda}\right)^4(N_{0c}|4a_0+a_c|^2+N_c|a_c|^2)\ \text{GeV}, \tag{12}$$

with $N_{0c}\approx3.46\times10^{-6}$ and $N_c\approx10.31\times10^{-6}$. In addition, if lepton universality is assumed, Eq. (12) is to be multiplied by 3 to account for all of the known neutrino species. From the LEP-2 data, the L3 Collaboration set the following limit on $Z\to\bar{\nu}\nu\gamma\gamma$ [7]:

$$\text{BR}(Z\to\bar{\nu}\nu\gamma\gamma)\leqslant3.1\times10^{-6}. \tag{13}$$

Assuming that either a_0 or a_c is dominant we obtain the following bounds:

$$\frac{|a_0|}{\Lambda^2}\leqslant0.106\ \text{GeV}^{-2}\ \text{ if }\ a_0\gg a_c, \tag{14a}$$

$$\frac{|a_c|}{\Lambda^2}\leqslant0.215\ \text{GeV}^{-2}\ \text{ if }\ a_c\gg a_0, \tag{14b}$$

which are weaker than those obtained at LEP-2 from $Z\gamma\gamma$ and $W^+W^-\gamma$ production [5]. However, it is important to note that the bounds based on the latter processes do not agree, as pointed out in Ref. [8], which means that much work along these lines is still required. In general Eqs. (12) and (13) yield an allowed area in the a_0/Λ^2 versus a_c/Λ^2 plane, as depicted in Fig. 2.

We can also take a different approach and instead of bounding the $ZZ\gamma\gamma$ coupling, we may use the most stringent bounds on it to predict its contribution to the $Z\to\bar{\nu}\nu\gamma\gamma$ decay. From the most stringent experimental bound on the a_0/Λ^2 and a_c/Λ^2 coefficients, of the order of 10^{-2}-10^{-3} [5], it follows a limit on the contribution of the $ZZ\gamma\gamma$ coupling to the $Z\to\bar{\nu}\nu\gamma\gamma$ decay:

$$\mathrm{BR}(Z \rightarrow Z^* \gamma\gamma \rightarrow \bar{\nu}\nu\gamma\gamma) \lesssim 1 \times 10^{-12}. \tag{15}$$

This indirect bound is above than the one found for the contribution of the neutrino-one-photon interaction $\bar{\nu}\nu\gamma$, which is of the order of 10^{-14} [10].

IV. CONCLUDING REMARKS

In closing it is interesting to note that the $ZZ\gamma\gamma$ coupling also contributes to the rare decay $Z \rightarrow \bar{l}^+ l^- \gamma\gamma$, for which an experimental bound has already been set [8]. This rare decay mode also receives contributions from the quartic $Z\gamma\gamma\gamma$ coupling. Therefore it would be possible, in principle, to use this rare decay to bound such QNGBCs. As far as the $Z\gamma\gamma\gamma$ coupling is concerned, a tighter bound on it can be obtained from the three-body decay $Z \rightarrow \gamma\gamma\gamma$. The latter, together with the $e^+e^- \rightarrow \gamma\gamma\gamma$ reaction, have been studied within the effective Lagrangian approach [15] and we will not repeat the same analysis here. From the result presented in Ref. [15] a bound on the effective $Z\gamma\gamma\gamma$ vertex can be derived, which is of the same order of magnitude of the one obtained in this work for the $ZZ\gamma\gamma$ coupling.

Finally we would like to emphasize that the importance of studying QNGBCs is rooted in the fact that they have a different origin than TNGBCs, i.e. they are induced by effective operators that do not induce any TNGBC. In this work we have found a constraint on the $ZZ\gamma\gamma$ coupling from the $Z \rightarrow \bar{\nu}\nu\gamma\gamma$ decay mode, which is weaker than the bounds derived from the analysis of $Z\gamma\gamma$ and $W^+W^-\gamma$ production at LEP-2. This is explained from the fact that the decay $Z \rightarrow \bar{\nu}\nu\gamma\gamma$ has a suppression factor due to the virtual Z boson propagator, which suffers less suppression when the initial Z boson is allowed to be off-shell.

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