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Direct tests of quartic vector bosons couplings in $\gamma\gamma \to W^+W^-$ and $\gamma\gamma \to ZZ$: The use of polarized spectra

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Abstract

We study the effects of anomalous four bosons couplings in boson pair production in high energy $\gamma\gamma$ collisions. We consider two anomalous couplings which satisfy a global weak $SU(2)$ symmetry and a local $U(1)_{QED}$. We emphasize the importance of using polarization to disentangle the two anomalous operators and single out the $J_Z = 0$ signal which test the scalar sector of the electroweak theory.

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1 Introduction

Within the standard description of weak interactions, the photon emerges as a combination of the hypercharge gauge boson and the neutral component of the weak boson triplet. The latter part connects directly to the as yet untested genuine non-abelian nature of the Standard Model (SM). These gauge couplings involve not only the trilinear $WW\gamma$ couplings to which an extensive literature has been devoted[1] but also the quadri-linear couplings which have been until very recently ignored. This state of affairs is due principally to a phenomenological motivation spurred on by the study of WW production at LEP-II[2] and to a lesser extent by single W production at HERA[3] where only the trilinear couplings are involved. Nonetheless at higher energies quartic couplings can be investigated either in triple vector boson production in e^+e^- and pp machines or in boson fusion at these colliders[4]. High-energy photon colliders appear almost as a unique tool to critically test the electromagnetic quartic couplings.

The measurement of quartic couplings that will become accessible to experimentalists in the future will provide just as important a test of the gauge structure of the standard model as the measurement of the trilinear couplings. Furthermore the origin of the two types of self-couplings could be different. For example, trilinear couplings could be the result of integrating out heavy fields while quartic couplings could represent the exchange of some heavy states. In particular the effective contact interactions involving vector bosons could correspond to a scalar exchange. In that sense the quartic couplings can indirectly test the unknown scalar sector of the standard model. Moreover it should be pointed out that there are models and approaches that generate anomalous quartic couplings while not affecting (or only slightly) the trilinear part. For example, new structures in quartic couplings appear in massive Higgs-less Yang-Mills theories[5], in the non-linear realization of symmetry breaking[6] or in alternative models of symmetry breaking[7]. Although the most drastic effects involve the massive gauge fields, two photons-two massive gauge bosons are also affected. The latter being the most accessible experimentally. The scalar connection constitutes one of the strongest theoretical motivations for studying the presence of anomalous quartic couplings and gives the photon collider the "edge" over other colliders since purely bosonic scatterings can be investigated.

In the following analysis our strategy will be to assume that the triple gauge boson coupling is standard and study only anomalous quartic couplings. This will allow us to illustrate the effects of quartic couplings. In the event where anomalous trilinear couplings are also present they will presumably be seen in high energy e^+e^- machines first. It would then be a simple task to incorporate these anomalous couplings in our analysis[8].

2 Anomalous couplings

We are concerned with direct tests of the gauge and symmetry breaking structure of the standard model and the detection of possible deviations from it. We choose a phenomenological approach which does not refer to a specific model. The only condition imposed are a custodial global $SU(2)$ symmetry and of course a $U(1)_{em}$ gauge symmetry since we are considering operators involving photons.

The lowest dimension operators involving two photons are dim-6 operators. There are two such operators that are C and P conserving

$$\mathcal{L}_6^0 = -\frac{\pi\alpha}{4\Lambda^2} a_0\, F_{\alpha\beta}F^{\alpha\beta}(\vec{W}_\mu.\vec{W}^\mu) \quad \text{and} \quad \mathcal{L}_6^c = -\frac{\pi\alpha}{4\Lambda^2} a_c\, F_{\alpha\mu}F^{\alpha\nu}(\vec{W}^\mu.\vec{W}_\nu) \qquad (1)$$

where $\vec{W}_\mu$ is the $SU(2)_w$ triplet. In the physical basis we should replace

$$\frac{1}{2}(\vec{W}_\mu.\vec{W}_\nu) \longrightarrow W^+_\mu W^-_\nu + \frac{1}{2c_W^2} Z_\mu Z_\nu \qquad (2)$$

where $c_W \equiv cos\theta_W$ and Λ is a scale which we will set to M_W when extracting the bounds on a_0 and a_c. Note that the first operator ("neutral") can parametrize the exchange of a neutral scalar. These four-couplings operators are not related to the standardised parametrisation of the trilinear anomalous couplings (λ_γ in the usual notation) which also generate some anomalous quartic couplings. This operator is easily distinguished from $\mathcal{L}_6^{0,c}$ since it contributes to $e^+e^- \to W^+W^-$ and $e\gamma \to \nu W$

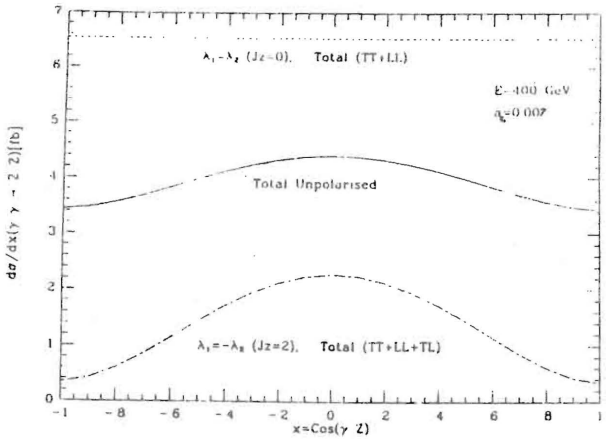

Figure 1: Angular distributions at $400GeV$ in $\sigma(\gamma\gamma \to ZZ)$ due to the a_c coupling for various polarizations

which involve only trilinear couplings[9]. Also the operator "λ" does not imply a $\gamma\gamma ZZ$ coupling in stark contrast with our case (see Eq. 2).

We should mention that presently there are virtually no constraints on our operators. Indeed, the $SU(2)$ global symmetry imposed on the Lagangian forbids any contribution to the ρ parameter. There is only a very weak constraint coming from $\Pi'_{\gamma\gamma}$ [10]. Furthermore the unitarity limits that could be set on quartic couplings from partial-wave analysis in $\gamma\gamma \to VV$ are orders of magnitude weaker than the direct ones we will study below.

The dimension-6 operators introduced above contribute to both $\gamma\gamma \to W^+W^-$ and $\gamma\gamma \to ZZ$. The helicity amplitudes with the anomalous quartic couplings can be written in a compact form.[11] For the "neutral" operator there is only a $J_Z = 0$ ($\lambda_1 = \lambda_2$) amplitude

$$\begin{aligned} \mathcal{A}^0_{\lambda_1\lambda_2;00} &= a_0^V(\lambda_1+\lambda_2)\ (\gamma-2) \quad \text{and} \quad \mathcal{A}^0_{\lambda_1\lambda_2;\lambda_3\lambda_4} = a_0^V\ (\lambda_1+\lambda_2)(\lambda_3+\lambda_4) \\ a_0^W &= \pi\alpha(M_W^2)\ \frac{s}{\Lambda^2}\ a_0 \quad , \quad a_0^Z = \frac{a_0^W}{c_W^2} \quad \text{and} \quad \gamma = \frac{s}{M_V^2} \end{aligned} \tag{3}$$

where λ_1, λ_2 are the helicities of the photons and $\lambda_3, \lambda_4 = \pm 1$ are the transverse Z polarizations. The vector bosons produced here would be mostly longitudinal (see Fig. 1) and this effect becomes more important at high energies.

The "charged" operator receives contribution from both the $J_Z = 2$ and $J_Z = 0$ channels. The dominant helicity amplitude for the longitudinal vector bosons is

$$\mathcal{A}^c_{\lambda_1\lambda_2;00} = a_c^V\ \left[(1+\lambda_1\lambda_2)\ (2-\gamma) + \frac{\gamma}{2}\ (1-\lambda_1\lambda_2)\sin^2\theta\right]$$

$$a_c^W = \pi\alpha\ \frac{s}{8\Lambda^2}a_c \quad \text{and} \quad a_c^Z = \frac{a_c^W}{c_W^2} \tag{4}$$

Clearly the angular distribution for the two operators is very different, the neutral giving a flat distribution and the charged one leading to most of the events in the central region. These characteristics suggest two possible ways of differentiating the anomalous operators. If polarization is available the different contribution to the $J_Z = 2$ channel will do the job, otherwise a careful analysis of the angular distribution of the outgoing vector bosons is necessary. As seen in Fig.1, the $J_Z = 0$ channel is by far the most sensitive one with a cross-section nearly five times larger. Of course the new physics could contribute to both operators. An interesting case is the one where $a_c = -2a_0 = a_n$, if this happens the $J_Z = 0$ contribution cancels out and we obtain only very weak constraints.

3 Analysis

3.1 *Photon spectrum*

Since both $\gamma\gamma \to W^+W^-$and ZZ cross-sections increase with energy because of the presence of the quartic couplings, it is best to study these reaction at as high $\gamma\gamma$ invariant masses as available without a loss in luminosity. The high-energy photon collider where the energetic photons are obtained through backscattering of laser light off initial electron beams are then the ideal $\gamma\gamma$ sources. The energy spectrum of the photons in such a collider were discussed in previous talks[8,12,13], we use the same luminosity functions. To achieve a higher degree of monochromaticity of the spectrum, polarization is essential. In Fig. 2, it is shown that when polarizing both the electron beam and the laser, a much larger fraction of the photons have an energy near the maximum ($\tau = \sqrt{s_{\gamma\gamma}}/\sqrt{s_{e^+e^-}}$). This is best achieved with the configuration where the circular polarization of the laser (P_c) and the mean helicity of the electron ($\lambda/2$) are opposite, for the $\gamma\gamma$ collider this means $\lambda P_c = \lambda' P'_c = -1$ ($'$ are for the opposite arm of the photon collider). Fig.2 shows the case

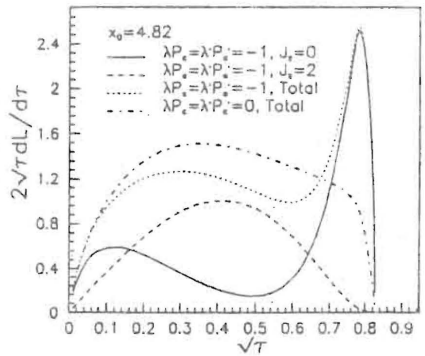

Figure 2: Spectrum of backscattered photons

where both lasers are tuned to have a right-handed circular polarization ($P_c = P'_c = +1$). This has the added advantage that the high-energy photons are produced mostly with the same helicity therefore giving a $J_Z = 0$ dominated environment. This channel can be isolated with a cut on the reduce energy $\sqrt{\tau} > 0.7$. The $J_Z = 2$ component can also easily be singled out, by flipping *both* the electron and laser polarization of *one* of the arms *only* while maintaining $\lambda P_c = -1$ (for a maximum of monochromaticity). In this case, the $J_Z = 0$ and $J_Z = 2$ spectrum of Fig. 2 will be interchanged. For some processes, the invariant mass of the final state may not be reconstructed if for instance weakly interacting neutrals are produced. We propose to still integrate over the whole spectrum by studying the following observable which measures the ratio of number of events (N) with opposite initial polarizations as

$$R = \frac{N(\lambda = -P_c = -1\ ;\ \lambda' = -P'_c = -1)}{N(\lambda = -P_c = 1\ ;\ \lambda' = -P'_c = -1)} \tag{5}$$

For the processes at hand a measurement of this ratio could distinguish between the two anomalous operators as will be shown in the following.

3.2 $\gamma\gamma \to ZZ$

In the SM, the amplitude for $\gamma\gamma \to ZZ$ is exactly zero at tree-level. If we impose C conjugation on the anomalous couplings then only only anommalous quartic $ZZ\gamma\gamma$ contribute to this reaction. Since global $SU(2)$ symmetry relates any deviation in $W^+W^-\gamma\gamma$ to that in $ZZ\gamma\gamma$, this reaction is the ideal place to search for the $SU(2)$ symmetric anomalous quartic couplings because of the absence of background from the SM. The most interesting aspect of this process is the Higgs production in the s-channel with its subsequent decay into a pair of Z's. One should remark that the a_0 operator is a good parametrization of this effect, apart from the resonant structure of the Higgs exchange. Hence our analysis of the a_0 operator can also serve as a guide for heavy Higgs ($M_H > 2M_Z$) studies in $\gamma\gamma$.

TABLE 1. *Limits on anomalous couplings from $\gamma\gamma \to ZZ$.*

	Ideal case $\sqrt{s_{\gamma\gamma}} = 400GeV$			Photon spectrum $\sqrt{s} = 500GeV$		$\sqrt{s} = 1TeV$
	unp.	$J_Z = 0$	$J_Z = 2$	unp.	pol.	pol.
$\lvert a_0\rvert <$	0.004	0.003		0.012	0.0067	0.0004
$\lvert a_c\rvert <$	0.007	0.006	0.01	0.02	0.018	0.0012

In order to put limits on the anomalous couplings, we will consider both the ideal case where all photons have an energy corresponding to 80% of the C.M. energy of the electrons and a more realistic case where the whole energy spectrum of the backscattered photons is included. In all instances we consider the possibility of polarizing the photons since it provides such an obvious way to distinguish the two operators a_0 and a_c. The best signal for Z pair production in $\gamma\gamma$ is the observation of one-sided events where one Z decays hadronically while the other decays into neutrinos. This signature corresponds to a combined ZZ branching ratio of 28%. A cut of $|\cos\theta| < 0.8$ while not sensibly reducing the event sample allows to totally get rid of two-(light)-fermions production, notably the two-(hard)-jets cross-section. We have required the observation of 10 one-sided events with the above topology and cut to set limits on the couplings. The luminosity was assumed to be $10fb^{-1}$. The results are given in Table 1 for both the ideal case and the initial photon spectrum. Because of the greater sensitivity of the $J_Z = 0$ channel, the bounds on a_0 are always smaller than those on a_c or a_n. For this last operator with the polarized spectrum at $500GeV$, we get $|a_n| < 0.032$. We also give the limits for $\sqrt{s} = 1TeV$ using a luminosity of $60fb^{-1}$ and requiring 20 events. The limits are an order of magnitude better than at lower energies.

In the two neutral gauge bosons production, the type of operator characterising the new physics is easily identifiable. The ratio R defined in Eq. 5 , is independent of the value of the anomalous operators. With the cuts given above and at $\sqrt{s} = 500GeV$, R is predicted to be 25.4, 2.5 and .08 for a_0, a_c and a_n respectively.

3.3 $\gamma\gamma \to W^+W^-$

In contrast with the neutral bosons production, the reaction $\gamma\gamma \to W^+W^-$ receives an important contribution from both the standard trilinear and quartic couplings. At $\sqrt{s_{\gamma\gamma}} = 400GeV$, the total cross-section is about $80pb$. Since most of the standard W's are produced close to the beam whereas the anomalous are more central (the purely anomalous contribution to the distribution is given by that of ZZ production in Fig. 1), a cut on the scattering angle eliminates a large fraction of standard model events thus enhancing possible signals from anomalous couplings. With a cut at 0.7, the cross-section is $17.6pb$ which still leaves a large number of events.

TABLE 2. *Limits on anomalous couplings from $\gamma\gamma \to WW$.*

Ideal case $\sqrt{s_{\gamma\gamma}} = 400GeV$	Unp.	$-0.06 < a_0 < 0.03$	$-0.10 < a_c < 0.035$
	$J_Z = 0$	$-0.05 < a_0 < 0.02$	$-0.10 < a_c < 0.036$
	$J_Z = 2$		$-0.02 < a_c < 0.02$
Photon spectrum $\sqrt{s_{e^+e^-}} = 500GeV$	Unp.	$-0.16 < a_0 < 0.055$	$-0.12 < a_c < 0.065$
	N^-	$-0.07 < a_0 < 0.03$	$-0.04 < a_c < 0.057$
	N^+		$-0.05 < a_c < 0.035$
	R	$-0.07 < a_0 < 0.03$	$-0.055 < a_c < 0.10$
Photon spectrum $\sqrt{s_{e^+e^-}} = 1TeV$	N^-	$-0.0035 < a_0 < 0.003$	$-0.007 < a_c < 0.0055$
	N^+		$-0.014 < a_c < 0.003$
	R	$-0.0035 < a_0 < 0.003$	$-0.004 < a_c < 0.006$

For this process the low statistics is not a problem, the main source of error comes

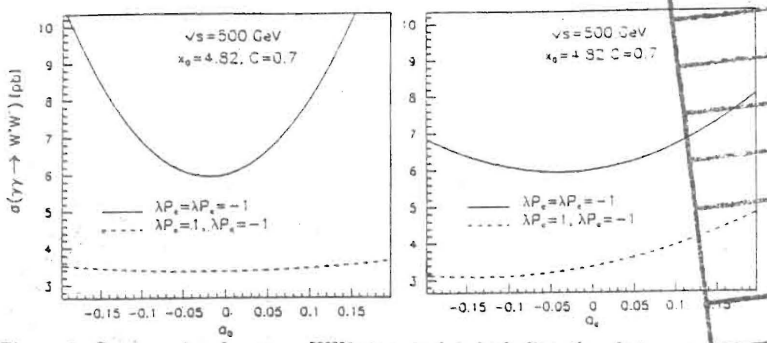

Figure 3: Cross-section for $\gamma\gamma \to WW$ vs a_0 and a_c including the photon spectrum.

from the systematics. To get limits on the anomalous couplings, we have assumed the same precision on the cross-section in the $J_Z = 0$, 2 and unpolarized case and taken $\frac{\Delta\sigma}{\sigma} = 3\%$ when considering the ideal spectrum case. When folding in the spectrum of the initial photons, we took $\frac{\Delta\sigma}{\sigma} = 5\%$ to account for possible additional uncertainties (like for example, simultaneaous measurements of all polarizations). The results are summarized in Table 2.

When including the photon spectrum a comparison between the limits obtained from the unpolarized cross-section or from the polarized one (see Fig. 3) clearly shows the advantage of using polarization. We define $N^-(N^+)$ as the number of events with the electron polarization $\lambda = -1(+1)$. These correspond to mostly the $J_Z = 0$ and $J_Z = 2$ channels respectively. Because of the interference with the standard model, the ratio R shown in Fig. 4, does not give a clear signal of the type of new physics as was the case in the ZZ mode. With the exception of large positive values, the two models are not easily distinguishable apart from the fact that N^+ is affected only by a_c. In Table 2 we give the bounds corresponding to a 5% measurement of R. Again by increasing the energy to $1TeV$, we get roughly an order of magnitude improvement on all bounds.

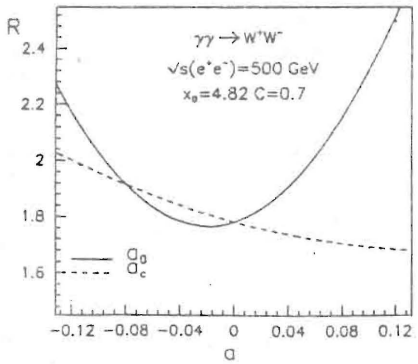

Figure 4: R vs anomalous couplings

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4 Conclusion

To interpret the limits just obtained, we can write $a_{0,c} = \alpha_H \frac{M_W^2}{\Lambda^2}$ where α_H is the strength of the new physics. If $\alpha_H \simeq 1$, as would be the case if the quartic coupling is a residual effect of strongly interacting W's at high scale, the limits obtained at $500GeV$ imply that one is sensitive to a scale Λ around $1TeV$. On the other hand for a weak coupling, $\alpha_H \simeq \alpha_W$, one is sensitive to a scale less than $200GeV$ in which case our parametrization for the exchange of a heavy scalar is obviously not valid. However a scalar with such a mass would probably be seen as a prominent resonance or/and produced directly in the e^+e^-mode.

To conclude, the process $\gamma\gamma \to ZZ$ offers the best test on quartic couplings. The limits we have obtained in Table 1 are almost two orders of magnitude better than the ones that could be obtained from $e^+e^- \to W^+W^-\gamma$ at the same energy[4]. Polarization is essential since it allows to isolate the $J_Z = 0$ and $J_Z = 2$ contributions, giving the possibility of disentangling a_0 from a_c. The process $\gamma\gamma \to W^+W^-$ does not give quite as good limits but can provide valuable information on the sign of a_0 and a_c.

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6 References

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