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Transverse beam polarization and CP-violating triple-gauge-boson couplings in $e^+e^- \to \gamma Z$

B. Ananthanarayan[a], Saurabh D. Rindani[b], Ritesh K. Singh[a], A. Bartl[c]

[a] *Centre for High Energy Physics, Indian Institute of Science Bangalore 560 012, India*

[b] *Theory Group, Physical Research Laboratory Navrangpura, Ahmedabad 380 009, India*

[c] *Universität Wien, Institut fur Theoretische Physik Boltzmanngasse 5, A-1090 Vienna, Austria*

Abstract

We show that an anomalous CP-violating $\gamma\gamma Z$ vertex gives rise to a novel asymmetry with transversely polarized electron and positron beams in the process $e^+e^- \to \gamma Z$. This asymmetry, which is odd under naive time reversal, is proportional to the real part of the $\gamma\gamma Z$ CP-violating coupling. This is in contrast to the simple forward-backward asymmetry of the γ (or Z) with unpolarized or longitudinally polarized beams studied earlier, which is even under naive time reversal, and is proportional to the imaginary part. We estimate the sensitivity of future experiments to the determination of CP-odd $\gamma\gamma Z$ and γZZ couplings using these asymmetries and transversely polarized beams.

1 Introduction

A future linear e^+e^- collider operating at a centre-of-mass (cm) energy of several hundred GeV would contribute greatly to a precise determination of the parameters of known particles and their interactions, as well as to the constraining of new physics. Longitudinal polarization of the e^+ and e^- beams, which is expected to be feasible at such colliders, would be helpful in reducing background as well as enhancing the sensitivity. It has been realized that spin rotators can be used to convert the longitudinal polarizations of the beams to transverse polarizations. The question has often been asked if such transverse polarization can be put to use to shed light on interactions or parameter ranges not accessible with longitudinal polarization, or to enhance their sensitivity. This question has not been discussed exhaustively in the current context as yet, though there have been some recent studies [1]-[7].

The role of transverse polarization in the context of CP violation has been studied in [6]-[10]. Since transverse beam polarization provides an additional reference coordinate axis in addition to the e^+e^- beam direction, there is the possibility of studying the azimuthal distribution of a single final-state particle. This has the advantage that the polarization of the produced particle, and hence its decay distribution, need not be measured. In [7] it was pointed out that an azimuthal distribution of a final-state particle A in a semi-inclusive process $e^+e^- \to A + X$ arising from the interference between a standard model (SM) contribution and a new-physics contribution arising at some high scale cannot contain a CP-violating part if the new-physics contribution arises from chirality-conserving vector (V) or axial-vector (A) type of interaction, neglecting the electron mass. This result, with the SM contribution restricted to a virtual photon exchange, can be deduced from the work of Dass and Ross [11]. In [7], this was generalized to include virtual Z exchange as well. On the other hand, chirality-violating scalar (S) and tensor (T) interactions can give rise to a simple CP-odd azimuthal asymmetry, as for example, in $e^+e^- \to t\bar{t}$ [7].

The above results were obtained with the condition that the SM contribution arises only through s-channel exchange of virtual photon and Z. The possibility of t- and u-channel exchange of an electron was not considered. Moreover, since the new physics is supposed to arise at a high scale, no t- or u-channel exchange of new particles was included. The results may get somewhat modified if these effects are taken into account. In particular, the t- or u-channel exchange would introduce an extra dependence on the scattering (polar) angle θ. In a process where A is its own conjugate, there may be a consequent forward-backward asymmetry corresponding to $\theta \to \pi - \theta$, which is CP odd. It is well-known that such an asymmetry could arise without transverse polarization (see, for example, [12, 13, 14]). However, such a forward-backward asymmetry, in the absence of transverse polarization, is even under naive time reversal T (*i.e.*, reversal of particle spins and momenta). Hence the CPT theorem implies that the contribution comes only from an absorptive part in one of the interfering amplitudes (see, for example, [15]). Thus, such a symmetry is only sensitive to the imaginary parts of the new-physics couplings.

In this paper we investigate the interesting possibility that if there is transverse polar-

ization, a T-odd but CP-even azimuthal asymmetry can be combined with the T-even but CP-odd forward-backward asymmetry to give an asymmetry which is both CP odd as well as T odd. In this case, the CPT theorem dictates that such an asymmetry meaure the real part of the new-physics couplings. The process we have chosen is $e^+e^- \to \gamma Z$, where the final-state particles are both self-conjugate*. This process occurs at tree level in SM. A CP-violating contribution can arise if anomalous CP-violating $\gamma\gamma Z$ and γZZ couplings are present. The interference of the contributions from these anomalous couplings with the SM contribution gives rise to the expected polar-angle forward-backward asymmetry, as well as new combinations of polar and azimuthal asymmetries. In particular, there is a CP-odd, T-odd asymmetry, which is proportional to the real part of the $\gamma\gamma Z$ coupling. This real part cannot be probed without transverse polarization†. There is an accidental cancellation of a similar contribution arising from real part of the γZZ coupling.

2 The process $e^+e^- \to \gamma Z$

We now describe the details of our work. The process considered is

$$e^-(p_-, s_-) + e^+(p_+, s_+) \to \gamma(k_1) + Z(k_2). \tag{1}$$

The most general effective CP-violating Lagrangian for $\gamma\gamma Z$ and γZZ interactions, consistent with Lorentz invariance and electromagnetic gauge invariance, and retaining terms upto dimension 6, can be written as

$$\begin{aligned} \mathcal{L} &= e\frac{\lambda_1}{2m_Z^2}F_{\mu\nu}\left(\partial^\mu Z^\lambda \partial_\lambda Z^\nu - \partial^\nu Z^\lambda \partial_\lambda Z^\mu\right) \\ &\quad + \frac{e}{16 c_W s_W}\frac{\lambda_2}{m_Z^2}F_{\mu\nu}F^{\nu\lambda}\left(\partial^\mu Z_\lambda + \partial_\lambda Z^\mu\right), \end{aligned} \tag{2}$$

where $c_W = \cos\theta_W$ and $s_W = \sin\theta_W$ and θ_W is the weak mixing angle. Terms involving divergences of the vector fields have been dropped from the Lagrangian as they would not contribute when the corresponding particle is on the mass shell, or is virtual, but coupled to a conserved fermionic current. Since we will neglect the electron mass, the corresponding current can be assumed to be conserved. We have not tried to impose full $SU(2)_L \times U(1)$ invariance, but only electromagnetic gauge invariance, as this is more general.

The SM diagrams contributing to the process (1) are shown in Figs. 1 (a) and 1 (b), which correspond to t– and a u–channel electron exchange, while the extra piece in the Lagrangian (2) introduces two s–channel diagrams with γ– and Z–exchange respectively, shown in Figs. 1 (c) and 1 (d). The corresponding matrix element is then given by

$$\mathcal{M} = \mathcal{M}_a + \mathcal{M}_b + \mathcal{M}_c + \mathcal{M}_d, \tag{3}$$

*A similar asymmetry has been considered for neutralino pair production in [6].

†An analogous situation is studied in [3], where transverse beam polarization allows one to probe certain CP-conserving triple gauge-boson couplings which cannot be probed with longitudinal polarization.

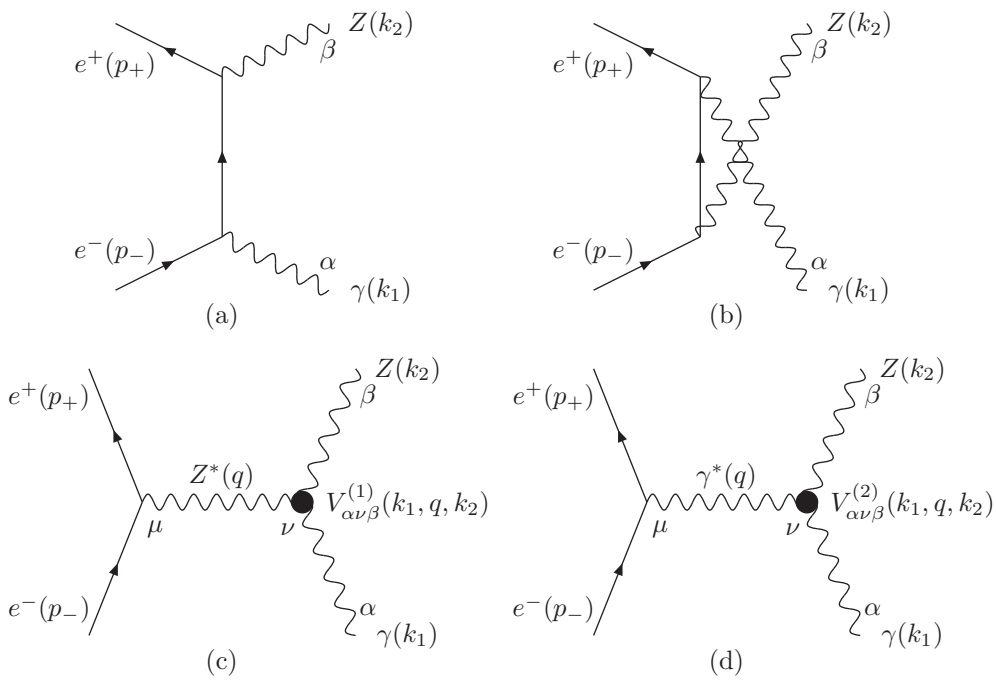

Figure 1: Diagrams contributing to the process $e^+e^- \to \gamma Z$. Diagrams (a) and (b) are SM contributions and diagrams (c) and (d) correspond to contributions from the anomalous γZZ and $\gamma\gamma Z$ couplings.

where

$$
\begin{aligned}
\mathcal{M}_a &= \frac{e^2}{4c_W s_W}\,\bar{v}(p_+)\,\not{\epsilon}(k_2)(g_V - g_A\gamma_5)\frac{1}{\not{p}_- - \not{k}_1}\,\not{\epsilon}(k_1)\,u(p_-),\\
\mathcal{M}_b &= \frac{e^2}{4c_W s_W}\,\bar{v}(p_+)\,\not{\epsilon}(k_1)\frac{1}{\not{p}_- - \not{k}_2}\not{\epsilon}(k_1)(g_V - g_A\gamma_5)u(p_-),\\
\mathcal{M}_c &= \frac{ie^2\lambda_1}{4c_W s_W m_Z^2}\,\bar{v}(p_+)\gamma_\mu(g_V - g_A\gamma_5)u(p_-)\frac{(-g^{\mu\nu} + q^\mu q^\nu/m_Z^2)}{q^2 - m_Z^2}V^{(1)}_{\alpha\nu\beta}(k_1,q,k_2)\epsilon^\alpha(k_1)\epsilon^\beta(k_2),\\
\mathcal{M}_d &= \frac{ie^2\lambda_2}{4c_W s_W m_Z^2}\,\bar{v}(p_+)\gamma_\mu u(p_-)\frac{(-g^{\mu\nu})}{q^2}V^{(2)}_{\alpha\nu\beta}(k_1,q,k_2)\epsilon^\alpha(k_1)\epsilon^\beta(k_2).
\end{aligned}
\tag{4}
$$

We have used $q = k_1 + k_2$, and the tensors $V^{(1)}$ and $V^{(2)}$ corresponding to the three-vector

vertices are given by

$$
\begin{aligned}
V^{(1)}_{\alpha\nu\beta}(k_1,q,k_2) &= k_1\cdot q\, g_{\alpha\beta}\, k_{2\nu} + k_1\cdot k_2 g_{\alpha\nu} q_\beta - k_{1\beta}\, q_\alpha\, k_{2\nu} - k_{1\nu}\, q_\beta\, k_{2\alpha} \\
V^{(2)}_{\alpha\nu\beta}(k_1,q,k_2) &= \frac{1}{2}\left[g_{\alpha\beta}\left(k_2\cdot q\, k_{1\nu} - k_1\cdot q\, k_{2\nu}\right) - g_{\nu\alpha}\left(k_2\cdot q\, k_{1\beta} + k_1\cdot k_2\, q_\beta\right)\right. \\
&\qquad \left. + g_{\nu\beta}\left(k_1\cdot k_2\, q_\alpha - k_1\cdot q\, k_{2\alpha}\right) + q_\alpha\, k_{2\nu}\, k_{1\beta} + q_\beta\, k_{1\nu}\, k_{2\alpha}\right].
\end{aligned}
\tag{5}
$$

In the above, the vector and axial vector Z couplings of the electron are given by

$$
g_V = -1 + 4\sin^2\theta_W; \quad g_A = -1. \tag{6}
$$

For compactness, we introduce the notation :

$$
\begin{aligned}
\overline{s} &\equiv \frac{s}{m_Z^2}\,, \\
\mathcal{B} &= \frac{\alpha^2}{16 s_W^2 m_W^2 \overline{s}}\left(1-\frac{1}{\overline{s}}\right)\left(g_V^2+g_A^2\right), \\
C_A &= \frac{\overline{s}-1}{4(g_V^2+g_A^2)}\left\{\left(g_V^2+g_A^2+(g_V^2-g_A^2)\,P_e P_{\overline{e}}\cos 2\phi\right)\mathrm{Im}\lambda_1\right. \\
&\qquad \left. -g_V(1+P_e P_{\overline{e}}\cos 2\phi)\,\mathrm{Im}\lambda_2 - g_A\, P_e P_{\overline{e}}\sin 2\phi\,\mathrm{Re}\lambda_2\right\}\,.
\end{aligned}
\tag{7}
$$

Using eqns. (3 - 7), we obtain the differential cross section for the process (1) to be

$$
\frac{d\sigma}{d\Omega} = \mathcal{B}\left[\frac{1}{\sin^2\theta}\left(1+\cos^2\theta+\frac{4\overline{s}}{(\overline{s}-1)^2} - P_e P_{\overline{e}}\frac{g_V^2-g_A^2}{g_V^2+g_A^2}\sin^2\theta\cos 2\phi\right) + C_A\cos\theta\right]\,, \tag{8}
$$

where θ is the angle between photon and the e^- directions, and ϕ is the azimuthal angle of the photon, with e^- direction chosen as the z axis and the direction of its transverse polarization chosen as the x axis. The e^+ polarization direction is chosen parallel to the e^- polarization direction. P_e and $P_{\overline{e}}$ are respectively the degrees of polarization of the e^- and e^+. We have kept only terms of leading order in the anomalous couplings, since they are expected to be small. The above expression may be obtained either by using standard trace techniques for Dirac spinors with a transverse spin four-vector, or by first calculating helicity amplitudes and then writing transverse polarization states in terms of helicity states [16].

We will assume a cut-off θ_0 on the polar angle θ of the photon in the forward and backward directions. This cut-off is needed to stay away from the beam pipe. It can further be chosen to optimize the sensitivity. The total cross section corresponding to the cut $\theta_0 < \theta < \pi - \theta_0$ can then be easily obtained by integrating the differential cross section above.

It is interesting to note that the contribution of the interference between the SM amplitude and the anomalous amplitude vanishes for $s = m_Z^2$. The reason for this is that for $s = m_Z^2$ the photon in the final state is produced with zero energy and momentum. As can be seen from eq. (5), the anomalous couplings vanish for $k_1 = 0$, leading to a vanishing intereference term.

In order to understand the CP properties of various terms in the differential cross section, we note the following relations:

$$\vec{P}\cdot\vec{k}_1 = \frac{\sqrt{s}}{2}|\vec{k}_1|\cos\theta\ , \tag{9}$$

$$(\vec{P}\times\vec{s}_-\cdot\vec{k}_1)(\vec{s}_+\cdot\vec{k}_1) + (\vec{P}\times\vec{s}_+\cdot\vec{k}_1)(\vec{s}_-\cdot\vec{k}_1) = \frac{\sqrt{s}}{2}|\vec{k}_1|^2\sin^2\theta\sin 2\phi\ , \tag{10}$$

$$(\vec{s}_-\cdot\vec{s}_+)(\vec{P}\cdot\vec{P}\vec{k}_1\cdot\vec{k}_1 - \vec{P}\cdot\vec{k}_1\vec{P}\cdot\vec{k}_1) - 2(\vec{P}\cdot\vec{P})(\vec{s}_-\cdot\vec{k}_1)(\vec{s}_+\cdot\vec{k}_1) = \frac{s}{4}|\vec{k}_1|^2\sin^2\theta\cos 2\phi\ . \tag{11}$$

where $\vec{P} = \frac{1}{2}(\vec{p}_- - \vec{p}_+)$. Observing that the vector $\vec{P}$ is C and P odd, that the photon momentum $\vec{k}_1$ is C even but P odd, and that the spin vectors $\vec{s}_\pm$ are P even, and go into each other under C, we can immediately check that only the left-hand side (lhs) of eq. (9) is CP odd, while the lhs of eqs. (10) and (11) are CP even. Of all the above, only the lhs of (10) is odd under naive time reversal T.

We now define the following CP-odd asymmetries, which combine a forward-backward asymmetry with an appropriate asymmetry in ϕ, so as to isolate appropriate anomalous couplings:

$$A_1 = \frac{1}{\sigma_0}\sum_{n=0}^{3}(-1)^n\left(\int_0^{\cos\theta_0} d\cos\theta\int_{\pi n/2}^{\pi(n+1)/2} d\phi\,\frac{d\sigma}{d\Omega} - \int_{-\cos\theta_0}^{0} d\cos\theta\int_{\pi n/2}^{\pi(n+1)/2} d\phi\,\frac{d\sigma}{d\Omega}\right) \tag{12}$$

$$A_2 = \frac{1}{\sigma_0}\sum_{n=0}^{3}(-1)^n\left(\int_0^{\cos\theta_0} d\cos\theta\int_{\pi(2n-1)/4}^{\pi(2n+1)/4} d\phi\,\frac{d\sigma}{d\Omega} - \int_{-\cos\theta_0}^{0} d\cos\theta\int_{\pi(2n-1)/4}^{\pi(2n+1)/4} d\phi\,\frac{d\sigma}{d\Omega}\right) \tag{13}$$

$$\begin{aligned} A_3 = \frac{2}{\sigma_0}\Bigg\{&\int_0^{\cos\theta_0} d\cos\theta\left(\int_{-\pi/4}^{\pi/4} d\phi\,\frac{d\sigma}{d\Omega} + \int_{3\pi/4}^{5\pi/4} d\phi\,\frac{d\sigma}{d\Omega}\right)\\ &-\int_{-\cos\theta_0}^{0} d\cos\theta\left(\int_{-\pi/4}^{\pi/4} d\phi\,\frac{d\sigma}{d\Omega} + \int_{3\pi/4}^{5\pi/4} d\phi\,\frac{d\sigma}{d\Omega}\right)\Bigg\} \end{aligned} \tag{14}$$

with

$$\sigma_0 \equiv \sigma_0(\theta_0) = \int_{-\cos\theta_0}^{\cos\theta_0} d\cos\theta\int_0^{2\pi} d\phi\,\frac{d\sigma}{d\Omega}\,. \tag{15}$$

These are easily evaluated to be

$$A_1(\theta_0) = -\mathcal{B}'\,g_A\,P_e P_{\bar{e}}\,\mathrm{Re}\lambda_2 \quad , \tag{16}$$

$$A_2(\theta_0) = \mathcal{B}'\,P_e P_{\bar{e}}\,((g_V^2 - g_A^2)\,\mathrm{Im}\lambda_1 - g_V\,\mathrm{Im}\lambda_2) \tag{17}$$

$$A_3(\theta_0) = \mathcal{B}' \left[\frac{\pi}{2}((g_V^2 + g_A^2)\,\mathrm{Im}\lambda_1 - g_V\,\mathrm{Im}\lambda_2) + P_e P_{\overline{e}}((g_V^2 - g_A^2)\,\mathrm{Im}\lambda_1 - g_V\,\mathrm{Im}\lambda_2) \right]. \tag{18}$$

$$\sigma_0 = 4\pi\mathcal{B} \left[\left\{ \frac{\overline{s}^2 + 1}{(\overline{s} - 1)^2} \ln\left(\frac{1 + \cos\theta_0}{1 - \cos\theta_0} \right) - \cos\theta_0 \right\} \right] . \tag{19}$$

In the above equations, we have defined

$$\mathcal{B}' = \frac{\mathcal{B}(\overline{s} - 1)\cos^2\theta_0}{(g_V^2 + g_A^2)\sigma_0(\theta_0)}. \tag{20}$$

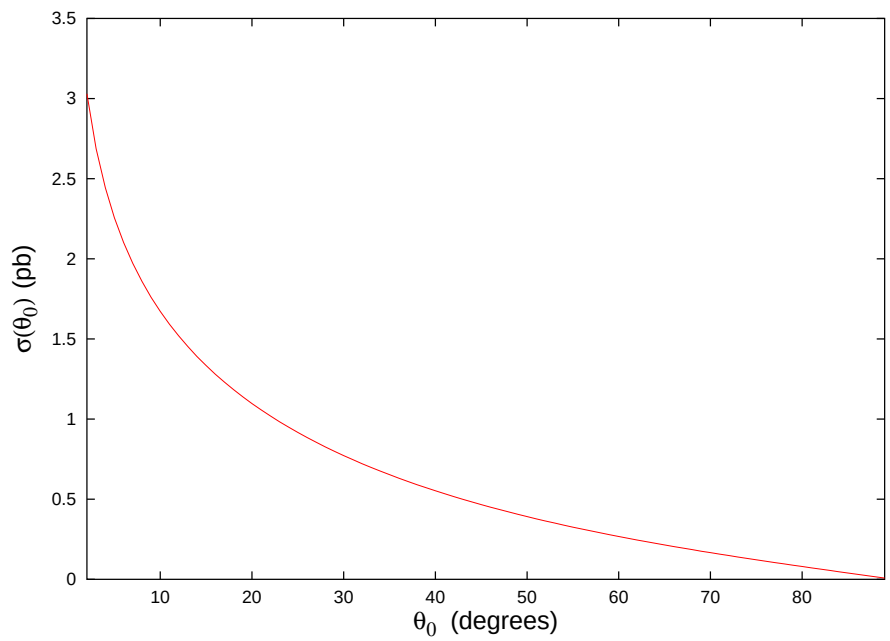

Figure 2: The SM cross section with a cut-off θ_0 in the forward and backward directions plotted as a function of θ_0.

We now make some observations on the above expressions which justify the choice of our asymmetries and highlight the novel features of our work. It can be seen that $A_1(\theta_0)$ is proportional to $\mathrm{Re}\lambda_2$, and the other two asymmetries depend on $\mathrm{Im}\lambda_1$ and $\mathrm{Im}\lambda_2$. Moreover, the latter two measured simultaneously can be used to get limits on the two couplings $\mathrm{Im}\lambda_1$ and $\mathrm{Im}\lambda_2$. It is interesting that A_1 does not depend on λ_1, which is the result of an accidental cancellation. This would not be the case, for example, if the Z in the s-channel exchange in Fig. 1(a) were different from the Z produced in the final state, so that their couplings to the the electron were different.

Note that the vector coupling g_V of the electron is small. As a result, the asymmetries A_2 and A_3 are relatively insensitive to $\mathrm{Im}\lambda_2$. However, in A_3, there is a partial cancellation of the $\mathrm{Im}\lambda_1$ contribution, making A_3 more sensitive to $\mathrm{Im}\lambda_2$ than A_2. This is borne out by our numerical results, see below.

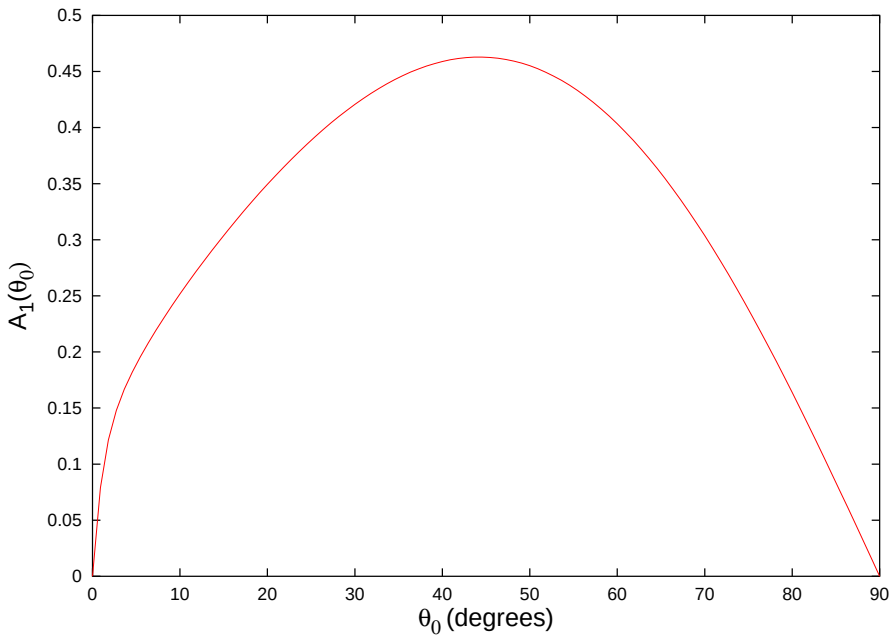

Figure 3: The asymmetry $A_1(\theta_0)$ defined in the text plotted as a function of the cut-off θ_0 for a value of Re $\lambda_2 = 1$.

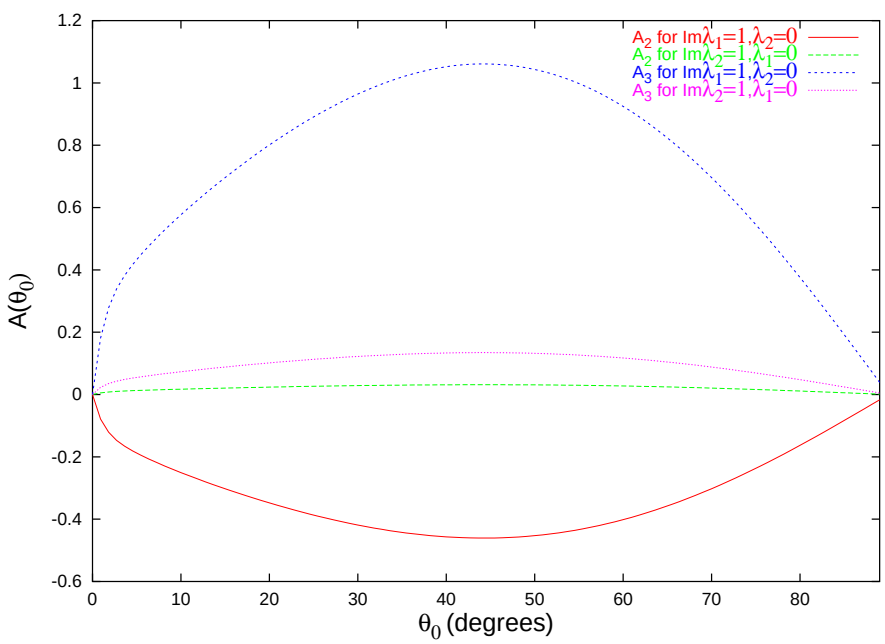

Figure 4: The asymmetries $A_2(\theta_0)$ and $A_3(\theta_0)$ defined in the text plotted as a function of the cut-off θ_0 for values Im $\lambda_1 = 1$, Im $\lambda_2 = 0$, and Im $\lambda_1 = 0$, Im $\lambda_2 = 1$.

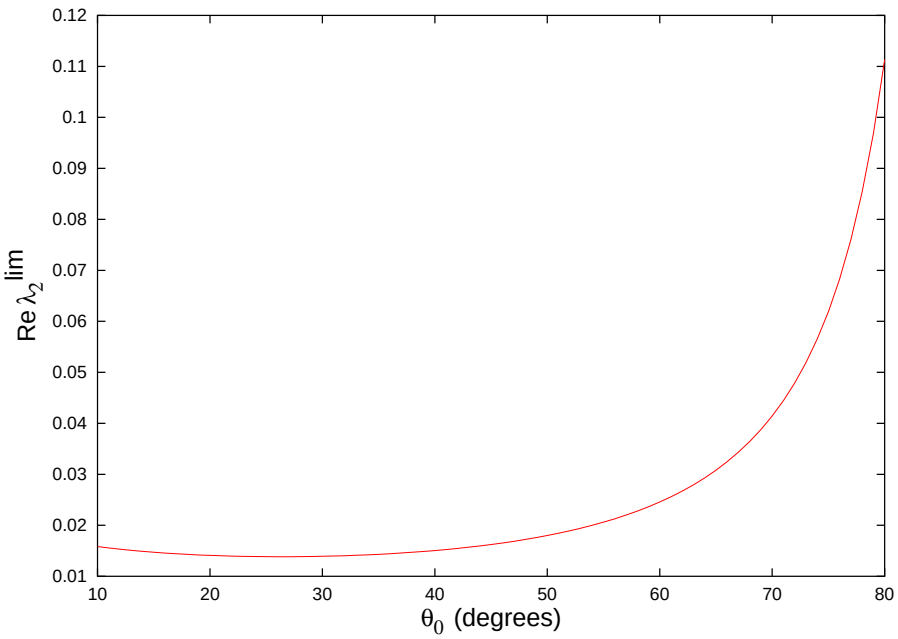

Figure 5: The 90% C.L. limit on Re λ_2 from the asymmetry $A_1(\theta_0)$ plotted as a funtion of the cut-off θ_0.

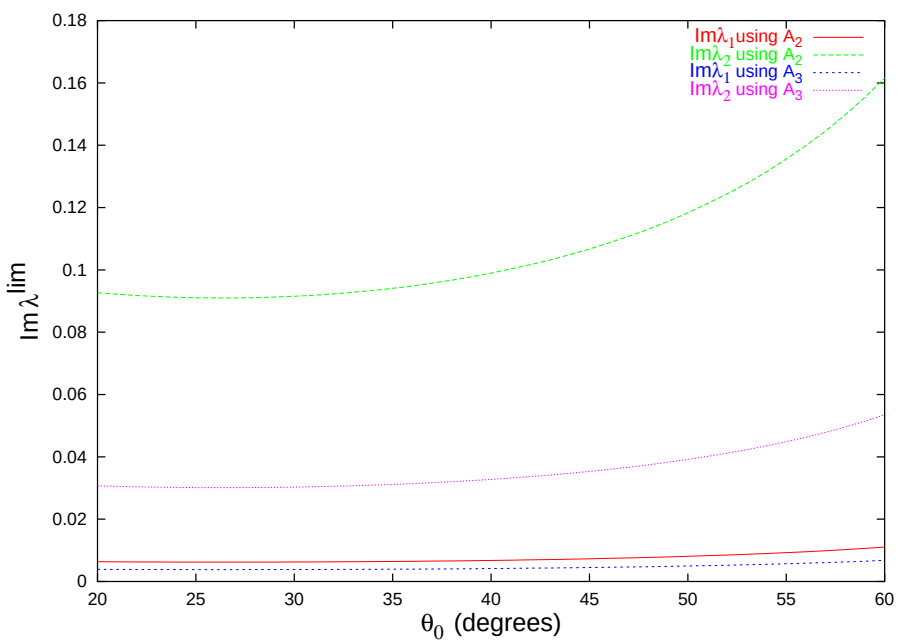

Figure 6: The 90% C.L. limits on Im λ_1 and Im λ_2, taken nonzero one at a time, from the asymmetries $A_2(\theta_0)$ and $A_3(\theta_0)$, plotted as a functions of the cut-off θ_0.

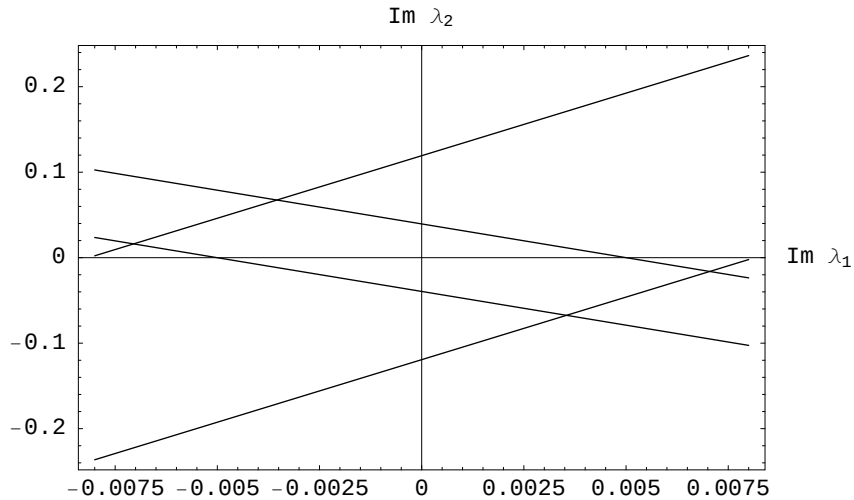

Figure 7: 90% CL contours for the simultaneous determination of Im λ_1 and Im λ_2. The region inside the trapezium is the allowed region.

3 Numerical Results

We now present our numerical results. The cross section with a cut-off θ_0 on θ is plotted in Fig. 2 as a function of θ_0. Figs. 3, 4 show the asymmetries as a function of the cut-off when the values of the anomalous couplings are taken to be nonzero one at a time. All the asymmetries vanish not only for $\theta_0 = 0$, by definition, but also for $\theta_0 = 90°$, because they are proportional to $\cos\theta_0$. They peak at around 45°.

We have calculated 90% CL limits that can be obtained with a linear collider with $\sqrt{s} = 500$ GeV, $\int Ldt = 500$ fb^{-1}, $P_e = 0.8$, and $P_{\overline{e}} = 0.6$ making use of the asymmetries A_i. The limiting value $\lambda^{\rm lim}$ (*i.e.* the respective real or imaginary part of the coupling) is related to the value A of the asymmetry for unit value of the coupling constant by

$$\lambda^{\rm lim} = \frac{1.64}{A\sqrt{N_{SM}}}, \tag{21}$$

where N_{SM} is the number of SM events.

A_1 depends on Re λ_2 alone, and can therefore place an independent limit on Re λ_2. We emphasize once again that information on Re λ_2 cannot be obtained without transverse polarization.

Fig. 5 shows the 90% CL limit on Re λ_2 as a function of the cut-off. The asymmetries A_2 and A_3 depend on both Im λ_1 and Im λ_2. Fig. 6 shows the 90% CL limits on Im λ_i taken to be nonzero one at a time, using the asymmetries A_2 and A_3. It can be seen from these figures that the limits are relatively insensitive to the cut-off at least for small values

Coupling	Individual limit from			Simultaneous limits
	A_1	A_2	A_3	
Re λ_2	1.38×10^{-2}			
Im λ_1		6.22×10^{-3}	3.82×10^{-3}	7.05×10^{-3}
Im λ_2		9.10×10^{-2}	3.01×10^{-2}	6.74×10^{-2}

Table 1: 90 % CL limits on the couplings from asymmetries A_i for a cut-off angle of 26°, $\sqrt{s} = 500$ GeV, and integrated luminosity of 500 fb^{-1}. The electron and positron transverse polarizations are assumed to be respectively 0.8 and 0.6.

of the cut-off. We find that the best limits are obtained for $\theta_0 = 26°$, though any nearby value of θ_0 would give very similar results. These correspond to Re $\lambda_2 = 0.0138$ (from A_1), Im $\lambda_1 = 0.00622$ (from A_2), Im $\lambda_1 = 0.00382$ (from A_3), Im $\lambda_2 = 0.0910$ (from A_2), and Im $\lambda_2 = 0.0301$ (from A_3).

As stated earlier, because of g_V being numerically small, the limits on Im λ_2, which appears in the expressions for the asymmetries multiplied by g_V, are worse than those on Im λ_1. However, it can also be seen that A_3 fares better than A_2 so far as Im λ_2 is concerned.

Finally, we have also evaluated the simultaneous 90% CL limits that can be obtained on Im λ_1 and Im λ_2 by measurement of A_2 and A_3. For this we have chosen $\theta_0 = 26°$. The corresponding contour for allowed values of the couplings for a null result of the measurement of A_2 and A_3 is shown in Fig. 7. This contour is obtained by equating the asymmetry obtained simultaneously from nonzero Im λ_1 as well as nonzero Im λ_2 to $2.15/\sqrt{N_{SM}}$. It can be seen that the simultaneous limits that can be obtained are weaker than individual limits, with numerical values Im $\lambda_1 = 0.00705$ and Im $\lambda_2 = 0.0674$. The best limits are summarized in Table 1.

4 Conclusions

In conclusion, we have studied a novel CP-violating asymmetry (A_1) in $e^+e^- \to \gamma Z$with anomalous neutral gauge boson couplings, which is special to a neutral final state, the observation of which needs both electron and positron transverse polarizations. This is of special interest in the context of the negative result stated in [7], that the observation of CP violation in a two-particle final state, without measuring the polarization of the final-state particles, is not possible with transversely polarized beams, unless there are chirality-violating couplings of the electron and positron. That result depended on an analysis where t- and u-channel particle exchanges were not taken into account.

Forward-backward asymmetry of a neutral particle with unpolarized or longitudinally polarized beams as a signal of CP violation has been studied before. However, the CPT theorem implies that in such a case the asymmetry is proportional to the absorptive part of the amplitude. The asymmetry A_1 that we study in the presence of transverse polarizations includes also an azimuthal angle asymmetry, which makes it odd under naive time reversal.

It is thus proportional to the real part of the anomalous coupling. This real part cannot be studied without transverse polarization.

We have also made a numerical study of the limits on various couplings that could be obtained at a future linear collider with $\sqrt{s} = 500$ GeV and an integrated luminosity of 500 fb^{-1} assuming realistic transverse polarizations of 80% and 60% respectively for e^- and e^+, respectively. The best limits are summarized in Table 1. We thus see that transverse polarization would provide a sensitive test of anomalous couplings, particularly, Re λ_2.

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