CP Nonconservation in Top Quark Production by (Un) Polarized $e^+e^-$ and $\gamma\gamma$ Collisions

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Abstract:
We report on an investigation of CP violation in (un)polarized $e^+e^- \rightarrow \bar{t}t$ resulting from an extended neutral Higgs sector or from the minimal SUSY extension of the Standard Model (SM). We consider c.m. energies from the $\bar{t}t$ threshold to the TeV range. In addition sensitivity estimates for CP-violating form factors of the top quark are made. Further we discuss the prospects of probing Higgs sector CP violation in $\gamma\gamma \rightarrow \bar{t}t$.

1 Introduction and Summary

In recent years quite a number of proposals have been made on how top quarks can serve as probes of CP-violating interactions beyond the Kobayashi-Maskawa (KM) mechanism ([1]-[17]). Model-independent analyses in terms of form factors can be found in [1, 2, 5, 7, 8, 10, 12]. Here we report on the results of an investigation of CP-violating interactions from an extended Higgs sector and from the minimal supersymmetric extension of the SM and their effects on $\bar{t}t$ production (and decay) at a linear collider. First we consider $e^+e^- \rightarrow \bar{t}t$. The new features of this study, as compared to previous work [1, 2, 5, 10, 12], are: we take into account the possibility of longitudinally polarized electron beams which enhance some of the effects, and we propose and study optimized observables with maximal sensitivity to CP effects for "semileptonic" $\bar{t}t$ events. For the Higgs model we find: (a) The highest sensitivities are reached somewhat above the $t\bar{t}$ threshold. (b) Longitudinal electron polarization would be an asset; yet the sensitivity of our best observable depends only weakly on the electron polarization. In the above models the effects in $e^+e^- \rightarrow t\bar{t}$ are due to non-resonant radiative corrections and are therefore not easy to detect. If a light Higgs

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boson $\varphi$ with mass $m_{\varphi} < 200$ GeV and sizable CP-violating couplings to top quarks exists then there is a chance to see a signal at a collider with 50 fb$^{-1}$ integrated luminosity. It would obviously be more promising to check for CP violation directly in $\varphi$ decays \cite{14, 17}. Effects in $t\bar{t}$ production due to a CP phase in gluino exchange are too small to produce statistically significant signals. 

In addition we have also investigated how well one could measure in the reactions (1), (2), (3) below the CP-violating form factors of the top quark with our optimized observables. Results are given in section 2. Of course, the sensitivity to the (dimensionful) form factors increases considerably with the c.m. energy and with the availability of beam polarization. (See also \cite{1, 12}.)

High energetic photon-photon collisions \cite{21}, which are discussed in the context of a linear collider, may provide, among other things, an interesting possibility to produce neutral Higgs bosons and study their quantum numbers. In the framework of two-Higgs doublet extensions we have investigated CP violation in unpolarized $\gamma\gamma \rightarrow \bar{t}t$ which includes resonant Higgs boson production and decay into $\bar{t}t$. Here the effects can be much larger than in $e^+e^-$-annihilation. If one or more Higgs bosons of intermediate mass $300$ GeV $\lesssim m_{\varphi} \lesssim 500$ GeV exist we find that $\gamma\gamma \rightarrow \varphi \rightarrow \bar{t}t$ would be a promising channel to study its/their CP properties. Sensitivity estimates will be given in section 3. For a detailed exposition of our studies and further references see \cite{6, 17}.

\section{\textbf{2} $e^+e^- \rightarrow \bar{t}t$}

In this section we consider the production of a top quark pair via the collision of an unpolarized positron beam and a longitudinally polarized electron beam:

$$e^+(e_+)+e^-(e_-,p) \rightarrow t(k_t)+\bar{t}(k_{\bar{t}}). \eqno(1)$$

Here $p$ is the longitudinal polarization of the electron beam ($p=1$ refers to right handed electrons). For our purposes the most interesting final states are those from semileptonic $t$ decay and non-leptonic $\bar{t}$ decay and vice versa:

$$t \bar{t} \rightarrow \ell^+(q_+)+\nu+\bar{b}+X_{\text{had}}(q_X), \quad \eqno(2)$$

$$t \bar{t} \rightarrow X_{\text{had}}(q_X)+\ell^-(q_-)+ \bar{\nu}+\bar{b}, \quad \eqno(3)$$

where the 3-momenta in eqs. (1) - (3) refer to the $e^+e^-$ c.m. frame. 

CP-violating interactions can affect the $\bar{t}t$ production and decay vertices. Quantum mechanical interference of the CP-even and -odd parts of the amplitudes for the above reactions then lead to the correlations which we are after. For CP-nonconserving neutral Higgs boson couplings and for CP-nonconserving gluino-quark-squark couplings it has been shown \cite{3} that these interactions lead to larger effects in $\bar{t}t$ production than in $t$ (and $\bar{t}$) decay. Therefore we consider observables which are predominantly sensitive to CP effects in the production amplitude which, in these SM extensions, arise at 1-loop through induced electric and weak dipole form factors $d_{\ell,Z}^t(s)$. The real parts $\text{Re}d_{\ell,Z}^t$ generate a difference in the $t$ and $\bar{t}$ polarizations orthogonal to the scattering plane of reaction (1), whereas non-zero absorptive parts $\text{Im}d_{\ell,Z}^t$ lead to a difference in the $t$ and $\bar{t}$ polarizations along the top direction of flight. The class of events (2), (3) is highly suited
to trace these spin-momentum correlations in the $\tilde{t}t$ production vertex through final state momentum correlations: From the hadronic momentum in (2), (3) one can reconstruct the $\tilde{t}$ and $t$ momentum, respectively and hence the rest frames of these quarks. The extremely short life time of the top quark implies that the top polarization is essentially undisturbed by hadronization effects and can be analyzed by its parity-violating weak decay $t \rightarrow b + W$, which we assume to be the dominant decay mode. Further, the charged lepton from semileptonic top decay is known to be by far the best analyzer of the top spin [15]. Therefore our observables are chosen to be functions of the directions of the hadronic system from top decay, of the charged lepton momentum, of the positron beam direction, and of the c.m. energy $\sqrt{s}$. In order to increase the statistical sensitivity of the observables we shall use the lepton unit momenta $\hat{q}_\pm^* \hbar$ in the corresponding top rest frames, which are directly accessible in the processes (2), (3).

The CP “asymmetries” which we discuss are differences of expectation values

$$A = \langle O_+(s, \hat{q}_+^*, \hat{q}_X^*, \hat{e}_+) \rangle - \langle O_-(s, \hat{q}_-^*, \hat{q}_X^*, \hat{e}_+) \rangle,$$

(4)

where the mean values refer to events (2), (3) respectively. For instance, $O_+ = (\hat{q}_+^* \times \hat{q}_X^*) \cdot \hat{e}_+$. The observable $O_-$ is defined to be the CP image of $O_+$. It is obtained from $O_+$ by the substitutions $\hat{q}_X^* \rightarrow -\hat{q}_X^*, \hat{q}_+^* \rightarrow -\hat{q}_-^*, \hat{e}_+ \rightarrow \hat{e}_-$. The ratio

$$r = \frac{\langle O_+ \rangle - \langle O_- \rangle}{\Delta O_+}$$

(5)

is a measure of the statistical sensitivity of $O$. Here $\Delta O = \sqrt{\langle O^2 \rangle - \langle O \rangle^2}$. For the observables used in this paper we have $\Delta O_+ \approx \sqrt{\langle O_+^2 \rangle}$ and $\Delta O_- \simeq \Delta O_-$. The signal-to-noise ratio of $A$ is given by $S_A = |r| \sqrt{N_{\text{event}}} / \sqrt{2}$, where $N_{\text{event}}$ is the number of events of type (2) or (3).

In [8] explicit expressions are given for two observables $O_\pm(1), O_\pm(2)$ with optimized sensitivity $|r|$, and these observables pick up dispersive ($\text{Re} d_{i}^{\gamma,Z}$) and absorptive ($\text{Im} d_{i}^{\gamma,Z}$) CP effects, respectively. In addition to the functional dependence exhibited in [8] they depend also on the electron beam polarization $p$. As the expressions are somewhat lengthy we do not reproduce them here. Schematically an optimized observable is constructed as follows: Consider a differential cross section which is of the form $d\sigma = d\sigma_0 + \lambda d\sigma_1$ where $\lambda$ is a small parameter. One can show [8, 14] that the observable with the highest statistical sensitivity is given by $O = d\sigma_1 / d\sigma_0$.

We shall consider phase space cuts which are CP-symmetric. When the $e^+e^-$ beams are unpolarized (or transversely polarized) the asymmetries [8] can be classified as being odd under a CP transformation. This means that contributions to $\langle O_\pm \rangle$ from CP-invariant interactions cancel in the difference. If the electron beam is longitudinally polarized the initial $e^+e^-$ state is no longer CP-symmetric in its c.m. frame and the CP classification no longer applies. Contributions from CP-conserving interactions can, in principle, contaminate $A$ if $p \neq 0$. However, in practice, this is not a problem because it can be argued that SM interactions induce contaminations at the per mill level [8]. On the other hand only ratios $|r| > 0.01$ have a chance to be detectable even at a high luminosity linear collider.

We now recall the salient features of neutral Higgs sector CP violation. For definiteness we consider two-doublet extensions of the SM with explicit CP violation in the Yukawa couplings (which leads to the Kobayashi-Maskawa phase) and in the Higgs potential. As a consequence the three physical neutral Higgs boson states $\varphi_{1,2,3}$ are in general states with
indefinite CP parity; i.e. they couple both to scalar and pseudoscalar quark and lepton currents with strength \( a_{jf}m_f/v \) and \( \tilde{a}_{jf}m_f/v \), respectively, where \( m_f \) is the fermion mass and \( v \simeq 246 \) GeV. For the top quark we have \( a_{jt} = d_2j/\sin\beta, \tilde{a}_{jt} = -d_3j\cot\beta \), where \( \tan\beta = v_2/v_1 \) is the ratio of the moduli of the vacuum expectation values of the two doublets, and \( d_2, d_3 \) are the matrix elements of a 3x3 orthogonal matrix which describes the mixing of the neutral Higgs states of definite CP parity. Only the CP=+1 components of the mass eigenstates \( \varphi_j \) couple to the \( W, Z, \) and charged Higgs bosons at the Born level. (For notation and details, see [3]).

CP violation requires that the neutral Higgs bosons are not mass-degenerate. Only if one of the bosons, say \( \varphi_1 \), is rather light then there is a chance that effects are detectable in \( e^+e^- \to \bar{t}t \). For the evaluation of the ratios \( r \) shown in Figs.1 we have put \( m_{\varphi_1} = 100 \) GeV and have assumed that \( m_{\varphi_{2,3}} >> m_{\varphi_1} \). Then the effect of \( \varphi_{2,3} \) on the dipole form factors is negligible. The EDM form factor of the top are proportional to \( \gamma_{CP} = -a_{1t}\tilde{a}_{1t} \) which is a measure of the strength of CP violation induced by \( \varphi_1 \) exchange. One may assume maximal CP violation in the neutral Higgs sector by putting \( d_{i1} = 1/\sqrt{3}(i = 1, 2, 3) \). The product \( \gamma_{CP} \) increases with decreasing \( \tan\beta \). Phenomenologically the experimental upper bounds on the neutron and electron EDMs give the rough upper bound \( |\gamma_{CP}| < 5 \). Nevertheless it would be interesting to obtain from the top system direct information on \( \gamma_{CP} \).

In Figs.1 we have plotted the ratios \( r_{1,2} \) for the optimized dispersive and absorptive observables \( O_{\pm}(i), O_{\pm}(2) \) defined in [6]. We have put \( m_t = 180 \) GeV, \( m_{\varphi_1} = 100 \) GeV, and \( \gamma_{CP} = 1 \). The ratios \( r_{1,2} \) are directly proportional to \( \gamma_{CP} \).

![Fig 1](image1.png)

Fig 1: Ratios \( r_1 \) (left figure) and \( r_2 \) (right figure) for the optimized dispersive and absorptive observables \( O_{\pm}(i), i = 1, 2 \) defined in [6] for \( m_t = 180 \) GeV, \( m_{\varphi_1} = 100 \) GeV, and \( \gamma_{CP} = 1 \).

Figs.1 show that the absorptive asymmetry has the highest sensitivity to \( \gamma_{CP} \) and that it depends only weakly on the beam polarization. The maximal sensitivity is reached at \( \sqrt{s} \simeq 450 \) GeV where \( |r_2| = 1.2\% (1\%) \) if \( p=\pm1 \) (0). In order to detect this as a 3 s.d. effect one would need 125000 (180000) events of the type \( \square \) and of \( \mathbb{E} \). These event numbers are unrealistically large. If the CP-violating effect is larger, say \( \gamma_{CP}=4 \), then only 1/16 of these events would be needed for a 3 s.d. signal. In order to reach a sensitivity \( \gtrsim 3 \) s.d. to couplings \( |\gamma_{CP}| \approx 4 \) an integrated luminosity of 50 fb\(^{-1} \) must be collected.
In general the asymmetries become smaller with increasing Higgs mass. If we keep $\gamma_{CP}$ fixed but change the Higgs mass to $m_\phi=200$ GeV then the maximal value at $\sqrt{s} \approx 450$ GeV of the ratio $r_2$ in Fig.1 is reduced by about 30%.

In summary, for light Higgs masses $m_\phi < 200$ GeV and sizable CP-violating coupling $\gamma_{CP}$ there is a chance to see Higgs sector CP violation as an induced dipole moment effect in $\bar{t}t$ production. A light Higgs particle $\phi$ would also be produced at a linear collider. A consequence of $\phi$ not being a CP eigenstate would be a CP violation effect in the $\phi$ fermion-antifermion amplitude at Born level which could be detected in $\phi \to \tau^+\tau^-$ [14]. (For further tests of the CP properties of Higgs particles, see [9, 15, 20].)

In the minimal supersymmetric extension of the SM additional CP-violating phases (besides the KM phase) can be present in the Majorana mass terms, e.g. of the gluinos, and in the squark (and slepton) mass matrices. For mass eigenstates these phases then appear for instance in the $\tilde{t}\bar{t}$–gluino couplings in the form (flavour mixing is ignored)

$$\mathcal{L}_{T\lambda\lambda} = i\sqrt{2} \ g_{QCD} \ \frac{e^{i\phi_t}}{\sin(2\phi_t)} \ \tilde{t}_R^a T^a (\bar{\lambda}_L^a t_L) + e^{-i\phi_t} \ \tilde{t}_L^a T^a (\bar{\lambda}_R^a t_R) + h.c. \quad (6)$$

with $\phi_t = \phi_\lambda - \phi_\tilde{t}$ and fields $\tilde{t}_{R,L}$ which are related to the fields $\tilde{t}_{1,2}$ corresponding to mass eigenstates by an orthogonal transformation. The 1-loop EDM and WDM form factors $d_{T,Z}$ generated by (6) are proportional to $g_{QCD}^2 \sin(2\phi_t)$. We have computed these form factors and the resulting asymmetries (4) for the optimized observables $O_{\pm}(1), O_{\pm}(2)$ for maximal CP violation, $\sin(2\phi_t) = 1$, and a range of SUSY masses $m_\lambda, m_{1,2} \geq 150$ GeV. We have found [6] that the ratios $|r_{1,2}| \leq 0.01$. Hence it is unlikely that a statistically significant signal can be detected at a 50 fb$^{-1}$ linear collider even for maximal SUSY CP violation. (Related studies were made in [3, 13].)

Furthermore we also investigated how well one could measure $Re d_{T,Z}$ and $Im d_{T,Z}$ independent of any model. For this purpose we computed the corresponding ratios $r$ for observables which, for a given polarization $p$, have optimized sensitivity to one of these form factors. Using an integrated luminosity of 20 fb$^{-1}$ (50 fb$^{-1}$) at $\sqrt{s} = 500$ GeV (800 GeV) we estimate the 1 s.d. statistical errors given in Table 1, assuming a tagging efficiency of 1 for the channels (2), (3).

|                  | 20 fb$^{-1}$, $\sqrt{s} = 500$ GeV | 50 fb$^{-1}$, $\sqrt{s} = 800$ GeV |
|------------------|-------------------------------------|-------------------------------------|
|                  | $p = 0$  | $p = +1$  | $p = -1$  | $p = 0$  | $p = +1$  | $p = -1$  |
| $\delta Re d_{T}$ | 4.6      | 0.86      | 0.55      | 1.7      | 0.35      | 0.23      |
| $\delta Re d_{Z}$ | 1.6      | 1.6       | 1.0       | 0.91     | 0.85      | 0.55      |
| $\delta Im d_{T}$ | 1.3      | 1.0       | 0.65      | 0.57     | 0.49      | 0.32      |
| $\delta Im d_{Z}$ | 7.3      | 2.0       | 1.3       | 4.0      | 0.89      | 0.58      |

Table 1: 1 s.d. sensitivities to the CP-violating form factors in units of $10^{-18}$ e cm.

Table 1 shows that the sensitivity to $Re d_{T}$ and to $Im d_{Z}$ grows substantially with beam polarization. For a fixed number of $t\bar{t}$ events the sensitivity to the nonrenormalizable “couplings” $d_{T,Z}$ obviously increases with increasing c.m. energy. However, one should keep in mind that the form factors vary with $\sqrt{s}$, and this variation is determined by thresholds which depend on the underlying dynamics of CP violation. If – like in the Higgs model – the relevant branch point where the form factors develop imaginary parts is set by the $t\bar{t}$ threshold then, as exemplified by Fig.1, there is no gain in going to very high energies. Needless to say: a priori one does not know.
3 $\gamma\gamma \rightarrow \bar{t}t$

After the discovery of a Higgs boson $\varphi$ at some collider a high energetic “Compton collider” of high luminosity could be tuned to $\gamma\gamma \rightarrow \varphi$ in order to perform a detailed study of the $\varphi$ quantum numbers (see,e.g.,[20]). If the photon polarizations are adjustable then one could check, as proposed in [16], for instance with the event asymmetry $(N(+,+) - N(-,-))/(N(+,-) + N(-,+))$ whether or not $\varphi$ is a CP eigenstate. (Here $\pm$ refer to the photon helicities.) For unpolarized $\gamma\gamma \rightarrow \varphi$ the CP property of $\varphi$ can be inferred from the final states into which it decays. If $\varphi$ is not a CP eigenstate then a CP-violating spin-spin correlation is induced in its fermionic decays already at the Born level, which can be as large as 0.5 [14]. It could be detectable in $\varphi \rightarrow \tau^+\tau^-$ and, for sufficiently heavy $\varphi$, in $\varphi \rightarrow \bar{t}t$. However, the narrow width approximation does not apply for a Higgs boson with mass in the vicinity or above the $\bar{t}t$ threshold and interference with the non-resonant $\gamma\gamma \rightarrow \bar{t}t$ amplitude decreases this spin-spin correlation significantly. In order to include also the case of light Higgs bosons $\varphi$ below the $\bar{t}t$ threshold we have computed, for two-Higgs doublet extensions of the SM, the complete set of CP-nonconserving contributions to $\gamma\gamma \rightarrow \bar{t}t$ in one-loop approximation [17]. In the spin density matrix for $\bar{t}t$ production by unpolarized photon beams this leads to CP-violating dispersive contributions with the spin structure $\hat{k}_t \cdot (s_+ \times s_-)$ and to absorptive contributions which are of the form $\hat{k}_t \cdot (s_+ - s_-)$ and to similar terms where $\hat{k}_t \rightarrow \hat{p}_\gamma$. Here $s_+,s_-$ are the spin operators of $t$ and $\bar{t}$, respectively. The latter structures correspond to polarization asymmetries.

As in the case of $e^+e^- \rightarrow \bar{t}t$ these spin correlations and polarization asymmetries are most efficiently traced through the ”semileptonic” $\bar{t}t$ decays (2) and the charge conjugated channels (3). The spin-spin correlation leads for events (2) to the triple product correlation $O_+ = \langle q_+ \times q_{\bar{t}t} \rangle \cdot \hat{q}_X$. Here the asterisk denotes the $\bar{t}$ rest system. The CP-reflected observable $O_-$ applies to the charge conjugated channels (3), and a non zero difference $\langle O_+ \rangle - \langle O_- \rangle$ would be an unambiguous sign of CP violation. A CP asymmetry with a somewhat higher sensitivity to Higgs boson induced effects is

$$A_{abs} = \langle q_+ \cdot \hat{q}_X \rangle - \langle q_- \cdot \hat{q}_X \rangle,$$

(7)

which is a consequence of the single spin asymmetries.

For the calculation of these asymmetries we used the normalized photon distributions $N(x)$ with energy fraction $x = E_{photon}/E_{beam}$ as given, e.g. in [21, 22]. The maximal energy fraction of a photon $x_{max}$ is determined by the laser energy $x_{max} = z/(1+z)$ with $z = 4E_{beam}E_{Laser}/m_{t^2}$. For a given beam energy the laser energy is chosen such that $z$ reaches its maximal value. This value is determined by the threshold for unwanted $e^+e^-$ pair production through annihilation of a backscattered photon with a laser photon. One gets $x_{max} \approx 0.8284$. Our result for the ratio $r$ based on the asymmetry (4) is shown in Fig.2 for $e^+e^-$ c.m. energy $\sqrt{s} = 500$ GeV and $m_t=180$ GeV as a function of the Higgs boson mass $m_\varphi$ for two sets of parameters: $d_{11} = 1/\sqrt{3}$, $\tan \beta = 0.627$ (set 1, solid line) and $d_{11} = 1/\sqrt{3}$, $\tan \beta = 0.3$ (set 2, dashed line). These two sets correspond to $\gamma_{CP}=1$ and $\gamma_{CP}=3.87$, respectively. In contrast to $e^+e^- \rightarrow \bar{t}t$ we have here a complicated dependence of the asymmetries on the Higgs boson couplings, because the diagrams in which the $\varphi$ propagator can become resonant contain also bosonic loops, and because in the resonant $\varphi$ propagator the total width $\Gamma_\varphi$ enters. The asymmetry $A_{abs}$ has an extremum when the Higgs mass is close to the $\bar{t}t$ threshold due to large interference effects and an additional extremum when
the Higgs mass is close to the maximal $\gamma \gamma$ energy. Even for a Higgs with mass above the maximal $\gamma \gamma$ energy there remains a significant interference effect that might be detectable.

![Graph showing the ratio $r$ for the asymmetry $A_{abs}$ as a function of the Higgs mass $m_\phi$ at $\sqrt{s} = 500$ GeV. The solid (dashed) line corresponds to the parameter set 1 (set 2).](image)

Fig 2: The ratio $r$ for the asymmetry $A_{abs}$ as a function of the Higgs mass $m_\phi$ at $\sqrt{s} = 500$ GeV. The solid (dashed) line corresponds to the parameter set 1 (set 2).

One may also determine, for a given Higgs mass, the $e^+e^-$ collider energy that maximizes the signal-to-noise ratio of an asymmetry. We have done this for $A_{abs}$ and the parameter set 1. The results are listed in Table 2. In brackets we also give the numbers for set 2. (Note that $S$ is not optimized with respect to this set.) In computing the number of events of the channels (2) or the charge conjugated channels (3) we use as in sect.1 semileptonic top decays into electrons, muons, and tauons.

| $m_\phi$ [GeV] | $\sqrt{s_{opt}}$ [GeV] | $N_{event}/(L/(100 \text{ fb}^{-1}))$ | $S/\sqrt{L}/(100 \text{ fb}^{-1})$ |
|---------------|------------------------|--------------------------------------|-----------------------------------|
| 100           | 700                    | $5.97(5.35) \times 10^3$             | 1.8(5.4)                          |
| 150           | 700                    | $5.96(5.31) \times 10^4$             | 1.8(5.3)                          |
| 200           | 700                    | $5.93(5.21) \times 10^4$             | 1.9(5.5)                          |
| 250           | 680                    | $5.44(4.70) \times 10^3$             | 2.1(6.1)                          |
| 300           | 650                    | $4.67(3.84) \times 10^3$             | 2.6(7.5)                          |
| 325           | 630                    | $4.10(3.30) \times 10^3$             | 3.0(8.9)                          |
| 350           | 570                    | $2.39(2.32) \times 10^4$             | 4.0(11.4)                         |
| 375           | 600                    | $3.41(4.27) \times 10^4$             | 4.4(8.8)                          |
| 400           | 710                    | $6.18(6.42) \times 10^4$             | 3.6(7.7)                          |
| 425           | 520                    | $1.87(2.31) \times 10^4$             | 3.7(4.7)                          |
| 450           | 560                    | $3.02(3.39) \times 10^3$             | 3.4(4.5)                          |
| 475           | 580                    | $3.58(3.85) \times 10^3$             | 3.3(4.2)                          |
| 500           | 610                    | $4.38(4.58) \times 10^3$             | 2.9(3.7)                          |

Table 2: Optimal collider energies, number of events of samples (2) or (3), and statistical significance $S$ of $A_{abs}$ for some Higgs masses and parameter set 1. The numbers in brackets are for set 2.
Table 2 tells us, for example, that if $m_\varphi = 350$ GeV (and if set 1 applies), one gets the largest sensitivity to the CP asymmetry $A_{abs}$ at a collider energy of 570 GeV. For an integrated $e^+e^-$ luminosity of 50 fb$^{-1}$ (and an $e\gamma$ conversion rate of one) we then get a statistical significance $S$ of 2.8(8.0). By using an optimized observable instead of (7) the sensitivity could still be enhanced. Should Higgs boson(s) of intermediate mass be discovered then, as our study shows, $\gamma\gamma \to \varphi \to \bar{t}t$ is a promising channel for probing Higgs sector CP violation.

References

[1] W. Bernreuther, O. Nachtmann, P. Overmann, T. Schröder: Nucl. Phys. B388 (1992) 53; ibid. B406 (1993) 516 (E)
[2] G.L. Kane, G. A. Ladinsky, C.–P. Yuan: Phys. Rev. D45 (1991) 124
[3] W. Bernreuther, T. Schröder, T.N. Pham: Phys. Lett. B279 (1992) 389
[4] C.R. Schmidt, M.E. Peskin: Phys. Rev. Lett. 69 (1992) 410
[5] W. Bernreuther, P. Overmann: Z. Phys. C61 (1994) 599
[6] W. Bernreuther, P. Overmann: preprint [hep-ph/9511250], to appear in Z. Phys. C
[7] T. Arens, L.M. Sehgal: Phys. Rev. D50 (1994) 4372
[8] D. Atwood, A. Soni: Phys. Rev. D45 (1992) 2405
[9] S. Bar-Shalom et al.: preprint SLAC-PUB-95-6981 (1995)
[10] J.P. Ma, A. Brandenburg: Z. Phys. C56 (1992) 97; A. Brandenburg, J.P. Ma: Phys. Lett. B298 (1993) 211
[11] B. Grzadkowski: Phys. Lett. B305 (1993) 384
[12] F. Cuypers, S.D. Rindani: Phys. Lett. B343 (1995) 333; S. Poulouse, S.D. Rindani: Phys. Lett. B349 (1995) 379; preprint [hep-ph/9509299]
[13] A. Bartl, E. Christova, W. Majerotto: Wien preprint HEPHY-PUB 624/95 (1995)
[14] W. Bernreuther, A. Brandenburg: Phys. Lett. B314 (1993) 104
[15] V. Barger, K. Cheung, A. Djouadi, B. Kniehl, P. Zerwas: Phys. Rev. D49 (1994) 79
[16] B. Grzadkowski, J.F. Gunion: Phys. Lett. B294 (1992) 361
[17] H. Anlauf, W. Bernreuther, A. Brandenburg: Phys. Rev. D52 (1995) 3803, Erratum to be published
[18] A. Czarnecki, M. Jezabek, J.H. Kühn: Nucl. Phys. B351 (1991) 70
[19] M. Diehl, O. Nachtmann: Z. Phys. C62 (1994) 397
[20] M. Kremer, J. Kühn, M. Stong, P. Zerwas: Z. Phys. C64 (1994) 21
[21] I.F. Ginzburg et al.: Nucl. Instrum. Meth. 205 (1983) 47; ibid. Vol. 219 (1984) 5; V.I. Telnov, Nucl. Instr. Meth. 294 (1990) 72
[22] J.H. Kühn, E. Mirkes, J. Steegborn: Z. Phys. C57 (1993) 615