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SPIN ASYMMETRY IN $\gamma e^+ \rightarrow e + 2 \text{ JETS}$:
A LABORATORY FOR FINAL STATE INTERACTIONS

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ABSTRACT

We isolate final state interactions not parametrized by fragmentation functions and yet present in QCD $O(\alpha_s)$ corrections to two-jet production in photon-longitudinally polarized electron scattering. Our calculation focuses on dynamical imaginary (absorptive) corrections that are absent in any tree diagrams. We explicitly demonstrate cancellation of imaginary infinities in the sum over graphs. The resulting asymmetry is predicted to be small. Such a quantity should be soon measured in experiments exploiting the naturally large electron polarization occurring in standard high energy storage rings.

I. INTRODUCTION

e^+e^- collisions are now recognized as one of the most powerful ways to probe the dynamics of the strong interaction. Among these reactions, the interest of the so-called photon-photon process is growing, as the increasing cross-section of the process makes it more and more suitable for a detailed analysis¹⁾. The proposal to use a laser beam in order to produce an intense and high energy photon beam from backward Compton scattering at the Stanford Linear Collider, together with the by now standard "single tagging" technique (used at PETRA and PEP and later at LEP) make the study of the photon-lepton collisions an even more promising field for the near future²⁾.

The process $\gamma\gamma^* \rightarrow \text{hadrons}$ can be studied at different levels of theoretical depth, depending on the information desired. The simple, inclusive process, $\gamma\gamma^* \rightarrow \text{anything}$, has been thoroughly analyzed and is useful as a testing ground of renormalization group controlled light-cone physics. Conversely, one may also study processes involving detected hadrons or jets -- $\gamma\gamma^* \rightarrow q\bar{q}$, $gg+X$ -- using parton models of jet fragmentation³⁾. This standard technique of modelling the final states, however, raises interesting questions about quark and gluon final state interactions. The purpose of this paper is to perturbatively study quark-antiquark final state interactions and present a novel experimental situation in which such interactions can be seen.

It is not easy to study final state interactions in QCD because in most processes their effects are simply parametrized into the fragmentation functions. In this paper we address the problem of single spin asymmetries in the production of two quark jets by a (quasi) real photon and a tagged, polarized lepton. Since electrons circulating in storage rings are known to naturally acquire a large polarization^{*)} which is efficiently transmitted to the virtual photon, there is a good probability that such an asymmetry can be measured⁴⁾. If final state interactions were solely confined to fragmentation, one could use symmetries of QCD to show that such an asymmetry is exactly zero -- regardless of the fragmentation function.

*) The natural polarization direction is transverse, perpendicular to the plane of the orbit. Standard techniques can be used to rotate it into longitudinal polarization required for our discussion.

The physical quantity which distinguishes non-trivial final state interactions from fragmentation processes is the presence of imaginary (absorptive) contributions from loop graphs in perturbation theory. Fragmentation type processes, no matter how complicated, yield a real contribution, hence a vanishing asymmetry. The lowest order non-vanishing contributions to the azimuthal asymmetry A , defined below, are given by the one-loop QCD diagrams of Fig. 1. In our calculation, A , a real number, is proportional to the imaginary parts of these amplitudes.

Despite the fact that A has been custom-made to pick out perturbative interactions not present in simple fragmentation, we still need fragmentation to describe further details of hadronization. However, fragmentation functions cannot be expected to absorb new infinities that may occur in a perturbative calculation of A . This is obvious because fragmentation is described by a real-valued probability, leaving no room for new imaginary infinities. In our calculation we find the infinities [terms of $O(\epsilon^{-1})$ in the dimensional regularization procedure] cancel in the sum over graphs. This is a non-standard and significant cancellation, necessary for us but not guaranteed within the usual applications of perturbative QCD.

We restrict ourselves to the subprocess $\gamma\gamma^* \rightarrow q\bar{q}$; the $O(\alpha_s^2)$ process $\gamma\gamma^* \rightarrow gg$ would lead to a vanishing azimuthal asymmetry since the two jets (which have opposite azimuth in the $\gamma\gamma$ centre-of-mass frame) are non-distinguishable.

We treat the outgoing quarks as on-shell free fermions, this approximation may be questioned. Indeed, treating them as slightly off-shell timelike objects would introduce new absorptive parts in the diagrams of the relevant order. This technically difficult problem is not addressed here^{*)}. The cancellation of divergences in our procedure is a signal (but not a proof) of the validity of this approximation.

We shall assume that the experiment is able to identify the present quark from the observation of a jet. This is expedited by the fact that the two-photon couplings very much suppress (by a factor 16) the creation of $d\bar{d}$ and $s\bar{s}$ quark-antiquark pairs, as compared to $u\bar{u}$ creation^{**)}. It is then

*) The contribution of such off-shell effects should be limited by the phase space available to the new thresholds, and this scales with the quark virtualness.

***) We neglect $c\bar{c}$ contributions. This is correct if we veto on the events with prompt K's.

sufficient, for instance, to identify the charge of the leading pion in the jet to assert with a good level of confidence whether we are indeed observing a u-quark (versus a $\bar{u}$ -antiquark) initiated jet. A more systematic study of the effect of jet fragmentation is quite straightforward and will not be addressed here.

The paper is organized as follows. Section II defines the kinematics of the reaction. In Section III, we present our calculation in some detail; an analytic expression of the asymmetry is given which we then analyze phenomenologically. Section IV is a brief conclusion. A table of imaginary parts of integrals evaluated in $d = 4-2\epsilon$ dimensions is relegated to an Appendix.

II. KINEMATICS

We consider the collision of a quasi-real photon γ with the spacelike virtual photon γ^* coming from a longitudinally polarized electron. In the electron-photon centre-of-mass frame, the four-vectors are defined as : $\lambda = (E, 0, 0, E)$ for the in-going electron; $\lambda' = (E', -E'\sin\theta, 0, E'\cos\theta)$ for the out-going electron; $q = (E-E', E'\sin\theta, 0, E-E'\cos\theta)$ for the virtual photon whose virtualness is $q^2 = -2EE'(1-\cos\theta)$. We will call $\sqrt{s}$ the total energy of the e- γ reaction, $\sqrt{s} = 2E$. In the case where the real photon is emitted by a positron, a boost along the positron momentum will lead to the over-all centre-of-mass frame.

Considering now the $\gamma\gamma^*$ centre-of-mass frame, the jet direction will be defined by the two angles (α, ϕ) in the following way (see Fig.2):

$$\begin{aligned} q &= ((W^2+q^2)/2W, 0, 0, (W^2-q^2)/2W) \text{ for the virtual photon,} \\ k &= ((W^2-q^2)/2W, 0, 0, -(W^2-q^2)/2W) \text{ for the real photon,} \\ p &= (W/2, (W\sin\alpha)/2, 0, (W\cos\alpha)/2) \text{ for the quark jet,} \\ r &= (W/2, -(W\sin\alpha)/2, 0, -(W\cos\alpha)/2) \text{ for the antiquark jet,} \\ \lambda &= (\lambda, \lambda\sin\theta'\cos\phi, \lambda\sin\theta'\sin\phi, \lambda\cos\theta') \text{ for the ingoing electron,} \\ &\text{where } W \text{ is the } \gamma\gamma^* \text{ total energy} \end{aligned}$$

$$W^2 = 4E(E-E') \tag{1}$$

and λ is the electron momentum

$$\ell = (s + q^2)/2W \quad (2)$$

The angle θ' is given in terms of s , W , and q^2 as

$$\cos \theta' = \frac{2sW^2}{(s+q^2)(W^2-q^2)} - 1 \quad (3)$$

α is thus the angle between $\gamma\gamma^*$ and the jet axis in the scattering plane (the xOz plane in the $\gamma\gamma^*$ centre-of-mass frame) while ϕ is the azimuthal angle between the bremsstrahlung plane (defined by the initial electron and the virtual photon, or the final electron) and the scattering plane. Note that to completely define ϕ we must distinguish the quark jet from the antiquark jet, i.e., the x axis from its opposite. Note also that this azimuthal angle cannot be defined when q^2 vanishes, i.e., when the electron emits a collinear photon. We will indeed find that the azimuthal asymmetry vanishes as $\sqrt{-q^2/s}$.

Let us conclude these remarks by listing some formulae that will be useful later

$$\begin{aligned} 4\varepsilon^{\mu\nu\rho\sigma} k_\mu p_\nu r_\rho \ell_\sigma &= -W(W^2-q^2) \ell \sin \theta' \sin \alpha \sin \phi \\ &= -W\sqrt{-q^2 s (s+q^2-W^2)} \sin \alpha \sin \phi \end{aligned} \quad (4)$$

$$\frac{d^3\vec{p}}{2p_0} \frac{d^3\vec{\ell}'}{2\ell'_0} = \frac{\pi}{8} |\vec{p}| s d|\vec{p}| d\cos\alpha dz' d\cos\theta d\phi \quad (5)$$

with $z' = 2E'/\sqrt{s}$. θ and z' define q^2 and W^2 as functions of s . Using the dimensionless character of the asymmetry will allow us to plot on the same graph results applicable either at PETRA or LEP energies. However, the luminosity factors, for instance the conversion rate from the positron beam to the "real" photon, are energy dependent.

III. CALCULATION

The cross-section we are interested in may be written as

$$\frac{2 p_0 2 \ell'_0 d\sigma(e^+ \gamma \rightarrow e q \bar{q})}{d^3 \vec{p} d^3 \vec{\ell}'} = \frac{(4\pi\alpha)^3}{(2\pi)^5} \frac{1}{2s} \int \frac{d^3 \vec{z}}{2z_0} \delta^4(\ell + k - \ell' - p - z) |M|^2 \quad (6)$$

where we have averaged over the real photon polarizations. $|M|^2$ can be decomposed into a leptonic part $L^{\mu\nu}$ and a hadronic part $W_{\mu\nu}$,

$$L^{\mu\nu}(\lambda) = \sum_{\lambda'} \bar{u}(\ell, \lambda') \gamma^\mu u(\ell, \lambda) \bar{u}(\ell, \lambda) \gamma^\nu u(\ell', \lambda') \quad (7)$$

where λ denotes the helicity of the incoming electron

$$\begin{aligned} W_{\mu\nu} &= \frac{C e_q^4}{2} \sum_{s, s', \alpha, \beta} \bar{u}(p, s) \Gamma_{\mu\alpha} \varepsilon^\alpha v(z, s') \bar{v}(z, s') \varepsilon^\beta \bar{\Gamma}_{\nu\beta} u(p, s) \\ &= -\frac{C e_q^4}{2} \sum_{s, s'} \bar{u}(p, s) \Gamma_{\mu\alpha} v(z, s') \bar{v}(z, s') \bar{\Gamma}_\nu^\alpha u(p, s) \end{aligned} \quad (8)$$

where C is the colour factor and e_q is the charge of the quark.

In the unpolarized case, we get at the Born term level

$$\frac{1}{2} [\mathcal{L}^{\mu\nu}(\lambda) + \mathcal{L}^{\mu\nu}(-\lambda)] = 2 \ell^\mu \ell'^\nu + 2 \ell'^\mu \ell^\nu + g^{\mu\nu} q^2 \quad (9)$$

while $C = N_c$, the number of colours, and

$$\Gamma_{\mu\alpha} = \Gamma_{\mu\alpha}^0 = \gamma_\mu \frac{1}{\cancel{p-q} + i\eta} \gamma_\alpha + \gamma_\alpha \frac{1}{\cancel{p-k} + i\eta} \gamma_\mu \quad (10)$$

where we neglect all fermion masses, leading to the result :

$$\sum_s |m|^2 = \frac{-16 N_c e_q^4 \Sigma}{q^2 (W^2 - q^2)^2 \sin^2 \alpha} \quad (11)$$

with

$$\begin{aligned} \Sigma = & 16 (p \cdot l^2 + r \cdot l^2) - 4 (p \cdot l + r \cdot l) (W^2 + q^2) \\ & + 4 (p \cdot l - r \cdot l) (W^2 - q^2) \cos \alpha + (q^2 + W^2)^2 \\ & + (W^2 - q^2)^2 \cos^2 \alpha, \end{aligned} \quad (12)$$

which is peaked at small q^2 and small transverse momentum.

The spin asymmetry $d\sigma(e^\uparrow \gamma) - d\sigma(e^\downarrow \gamma)$ involves an imaginary leptonic part

$$L_a^{\mu\nu} = -4i \varepsilon^{\alpha\beta\mu\nu} l_\alpha q_\beta \quad (13)$$

Simple i counting^{*)} shows that the Born approximation to $W_{\mu\nu}$ does not contribute to the spin asymmetry. It is easy to show that, among one-loop corrections, only those corresponding to graphs shown on Fig. 1 provide the needed imaginary part. Gluon emission diagrams or one-loop diagrams corresponding to vertex corrections without timelike legs, (Fig.3), do not yield any imaginary part. We thus calculate $W_{\mu\nu}$ at the lowest non-trivial order as the interference of the Born amplitude with the $O(\alpha_s)$ one-loop amplitude of Fig. 1. We obtain

*) This argument is metric-dependent. The absorptive nature of the contribution we study is not.

$$\text{Im } \Gamma_{\mu\alpha}^1 = \text{Im} \int \mu^{2\epsilon} \frac{d^n m}{(2\pi)^n} \left[\gamma_\rho \frac{1}{p-m+i\eta} \gamma_\mu \frac{1}{p-q-m+i\eta} \gamma_\alpha \frac{1}{-r-m+i\eta} \gamma_{\rho'} \right. \\ \left. + \gamma_\rho \frac{1}{p-m+i\eta} \gamma_\alpha \frac{1}{p-k-m+i\eta} \gamma_\mu \frac{1}{-r-m+i\eta} \gamma_{\rho'} \right] \frac{-g^{\rho\rho'}}{m^2+i\eta} . \quad (14)$$

The colour factor in that case is easily calculated to be $N_C C_F$, with $C_F = 4/3$ in QCD. The trace has been evaluated with the help of the symbolic program Reduce⁵⁾ while the imaginary parts of loop integrals are calculated using a dimensional regularization procedure, setting $n = 4-2\epsilon$. We tabulate in the Appendix the values of the needed integrals. While many integrals in fact diverge when $\epsilon \rightarrow 0$, each ϵ^{-1} being attached to a $\ln(-s/\mu^2)$, the complete result shows a cancellation of these divergences. This is in itself a non-trivial result ensuring the applicability of perturbative QCD to the quantity we have defined. We get :

$$L_a^{\mu\nu} W_{\mu\nu}^1 = \frac{4e_q^4 \alpha_s C_F N_C \epsilon^{\alpha\beta\gamma\delta} k_\alpha p_\beta r_\gamma \ell_\delta}{q^2 k.p^2 k.r^2 (W^2-q^2)^2} . \quad (15)$$

$$\left[A \ln \frac{W^2-q^2}{2k.p} + B \ln \frac{W^2-q^2}{2k.r} + C \right]$$

with

$$A = W^2 (W^2-q^2)^2 k.p \quad (16)$$

$$B = -W^2 (W^2-q^2)^2 k.r$$

$$C = 4 (k.p - k.r) q^2 k.p k.r$$

It may be noted that the bracket is antisymmetric with respect to p and r , so that $L_a^{\mu\nu} W_{\mu\nu}^1$ is symmetric but vanishes when $kp = kr$, i.e., when the two jets are emitted at 90° in the centre-of-mass system of the two photons. This is a consequence of the charge conjugation plus one initial state of two photons. We may rewrite (15) as

$$L_a^{\mu\nu} W_{\mu\nu}^1 = \frac{-64 e_q^4 \alpha_s C_F N_c W^3 \sqrt{s(s+q^2-W^2)}}{\sqrt{-q^2} (W^2 - q^2)^3} \cdot \frac{\sin \phi}{\sin \alpha} \left\{ \frac{\ln(1+\cos \alpha)/2}{1-\cos \alpha} - \frac{\ln(1-\cos \alpha)/2}{1+\cos \alpha} - \frac{q^2 \cos \alpha}{2W^2} \right\}^{(17)}$$

Compared to the unpolarized case, this expression displays some remarkable features. At low q^2 , it diverges only as $1/\sqrt{-q^2}$, leading to a vanishing asymmetry. At low α , it behaves like $1/\sin \alpha$ (the bracket remaining finite) to be compared with $1/\sin^2 \alpha$ in the unpolarized case. Its azimuthal dependence is very characteristic of a helicity asymmetry.

Let us now define first a differential asymmetry $a(\vec{p}, \vec{\ell}')$ as

$$a(\vec{p}, \vec{\ell}') = \frac{d\Delta\sigma(e\gamma \rightarrow eq\bar{q})/d^3\vec{p} d^3\vec{\ell}'}{d\bar{\sigma}(e\gamma \rightarrow eq\bar{q})/d^3\vec{p} d^3\vec{\ell}'} \quad (18)$$

with

$$\Delta\sigma = \sigma(e^\uparrow \gamma \rightarrow eq\bar{q}) - \sigma(e^\downarrow \gamma \rightarrow eq\bar{q})$$

and the denominator is the average over initial helicities. In all what follows, we will neglect the momentum dependence of the running coupling constant. To decide what is the relevant argument in $\alpha_s(Q^2)$ is not a question one can answer at such a lowest-order level. We will simply take $\alpha_s = 0.25$, neglecting the Q^2 evolution. The differential asymmetry (18) is a function of the angles ϕ , α and Θ and of the ratio $z' = 2E'/\sqrt{s}$ which is related to the value of $-q^2$. The ϕ dependence is particularly interesting: apart from a weak $\cos \phi$ dependence in the denominator, it is mostly due to the $\sin \phi$ factor of the numerator. We show in Fig. 4 this ϕ dependence for some typical values $\alpha = 45^\circ$, $\Theta = 50^\circ$, z' varying between 0.07 and 0.35. The magnitude of the effect can be read off this graph: it is of the order of a few per cent.

We can take advantage of this particular ϕ dependence to isolate the signal we are interested in by defining an integrated left hemisphere versus right hemisphere asymmetry:

$$A = \frac{\left[\int_0^\pi d\phi - \int_{-\pi}^0 d\phi \right] d\Delta\sigma / d\phi dz' d\cos\theta d\cos\alpha}{\int_0^{2\pi} d\phi d\bar{\sigma} / d\phi dz' d\cos\theta d\cos\alpha} \quad (19)$$

Note that the $\cos\phi$ dependence of the spin average case would not allow any contribution to such a left/right asymmetry. Using the scale invariance -- i.e., energy independence when $\sqrt{-q^2/s}$ is fixed -- of this quantity (α_s is taken constant), we present on Fig. 5a our results as a function of Θ and $\sqrt{-q^2/s}$ at $\alpha = 45^\circ$. They can thus be used for high energy existing machines as well as for very high energy future experiments. In order to discuss the detectability of the signal, we should worry about the expected counting rate in the region where A is not too small. This is shown on Fig. 5b, where the unpolarized cross-section [the denominator of Eq. (19)] is given in units of the $e^+e^- \rightarrow \mu^+\mu^-$ cross-section $\sigma_0 = 4\pi\alpha^2/3s$.

The conclusion of this analysis is quite straightforward. Perturbative QCD predicts a small single spin asymmetry whenever the counting rate is not extremely small. This is partly due to the kinematic $\sqrt{-q^2/s}$ factor attached to the definition of an azimuthal asymmetry, the counting rate being driven by a $1/q^2$ fall-off. A high total luminosity experiment at high energy may, however, be able to detect the predicted signal by taking advantage of all the features of the asymmetry we have presented.

IV. CONCLUSIONS

The purpose of this analysis has been to estimate certain final state interactions that cannot be parametrized into fragmentation functions. We have found that one-loop perturbative QCD interactions, in fact, are quite small in most kinematic regions. The natural question at the end of any perturbative calculation is the eventual importance of higher order corrections. Although we have no answer to this point, it may be appropriate to recall that spin asymmetries have been shown to be quite

unaffected by higher order corrections in other cases⁶⁾. This may be understood by the observation that an important part of the QCD radiative corrections may be factorized and effectively disappear in ratios such as asymmetries. On the other hand, one can adopt the position that the perturbative analysis only applies to the momentum dependence of the strong interaction phases⁷⁾ involved in A. Not enough is known from the technical development of QCD to decide between these alternatives at present.

Moreover, very little experimental work has been done to measure interference quantities (such as A) at high energies. QCD is not that well established that we can dispense ourselves of the effort of testing its spin predictions⁸⁾. Lepton beams in e^+e^- machines acquire a natural polarization which indeed should be used to get as much information as possible on the structure of final state interactions. This situation is complementary to the one in the Drell-Yan process⁹⁾, where there is a mixture of initial and final state interactions, and quantities involving interference can be measured. In either case there is comparatively more to be learnt from these experiments: the results would be of theoretical interest whatever the outcome.

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A P P E N D I X

We tabulate here the imaginary part of the integrals used in Section III. We use a dimensional regularization procedure by setting the space-time dimension to $n = 4 - 2\epsilon$. We define for any polynomial P

$$\frac{i\pi}{16\pi^2} \mathcal{I}[P] \equiv \Im m \frac{1}{i} \int \mu^{4-n} \frac{d^n m}{(2\pi)^n} \frac{P}{D} \quad (\text{A.1})$$

where

$$D = (m^2 + i\eta) \cdot ((p-m)^2 + i\eta) \cdot ((m+r)^2 + i\eta) \cdot ((p-m-k)^2 + i\eta) .$$

It turns out to be efficient to separate the divergent part (when $\epsilon \rightarrow 0$) under the form

$$\Delta = \epsilon^{-1} - \ln \frac{s}{\mu^2} + \ln 4\pi - \gamma_E$$

where γ_E is the Euler constant and $s = 2p \cdot r$. We first note the identities

$$\begin{aligned} \mathcal{I}[(p-m)^2 \cdot K] &= 0 \\ \mathcal{I}[(m+r)^2 \cdot K] &= 0 \end{aligned} \quad (\text{A.2})$$

for any polynomial K . They follow from the absence of physical threshold of three point Green functions with external legs that are all spacelike or lightlike. This allows us to express all the integrals needed in Section III in terms of the few listed below using relations systematically deduced from (A.1), e.g.,

$$\mathcal{I}[p \cdot m] = \frac{1}{2} \mathcal{I}[m^2] = -\mathcal{I}[r \cdot m]$$

The remaining integrals are straightforwardly calculated using $\text{Im } \ln(-s - i\eta) = -i\pi\theta(s)$:

$$\mathcal{I}[1] = \frac{-\Delta}{k \cdot p \ s} - \frac{\ln[(s-q^2)/2k \cdot p]}{k \cdot p \ s}$$

$$\mathcal{I}[m^2] = \frac{\Delta}{s - q^2}$$

$$\begin{aligned}
 I[k.m] &= \frac{-\Delta}{2s} - \frac{\ln[(s-q^2)/2k.p]}{s} \\
 I[k.m^2] &= \frac{-\Delta k.p}{2s} + \frac{k.r - k.p}{2s} - \frac{k.p \ln[(s-q^2)/2k.p]}{s} \\
 I[k.m^3] &= \frac{-\Delta k.p^2}{2s} - \frac{3k.p^2 + k.r^2 - 2k.p k.r}{4s} - \frac{k.p^3 \ln[(s-q^2)/2k.p]}{s} \\
 I[m^2 p.m] &= \frac{-\Delta s k.p}{(s-q^2)^2} - \frac{s}{2(s-q^2)} + \frac{2k.r s}{(s-q^2)^2} \\
 I[m^2 k.m] &= \frac{\Delta k.p}{s-q^2} + \frac{1}{2} \\
 I[m^2 p.m^2] &= \frac{\Delta k.p^2 s^2}{(s-q^2)^3} - \frac{s^2}{8(s-q^2)} - \frac{k.p s^2}{2(s-q^2)^2} \\
 &\quad + \frac{3k.p^2 s^2}{(s-q^2)^3} \\
 I[m^2 k.m^2] &= \frac{\Delta k.p^2}{s-q^2} + \frac{3k.p}{4} - \frac{k.r}{4} \\
 I[m^2 p.m k.m] &= -\frac{\Delta k.p^2 s}{(s-q^2)^2} - \frac{s}{4} - \frac{k.p s}{4(s-q^2)} \\
 &\quad + \frac{k.p k.r s}{(s-q^2)^2}
 \end{aligned}$$

where we have neglected $O(\epsilon)$ terms.

These results are of course specific to the case we study here, namely $s > 0$ and $q^2 < 0$, the quark legs being strictly on shell ($p^2 = r^2 = 0$).

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FIGURE CAPTIONS

Figure 1 : $O(\alpha_s)$ final state interactions contributing to single spin asymmetries in $e\gamma$ collisions.

Figure 2 : Kinematics of the reaction in the $\gamma^* \gamma$ centre-of-mass frame. $\vec{p}$ and $\vec{r}$ are in the xz plane.

Figure 3 : $O(\alpha_s)$ corrections which do not contribute to the spin asymmetry.

Figure 4 : The differential spin asymmetry at $\alpha = 45^\circ$ and $\theta = 50^\circ$ for different values of $z' = 2E'/\sqrt{s}$.

Figure 5 : a) The integrated asymmetry at $\alpha = 45^\circ$ as a function of θ and $\sqrt{-q^2/s}$. For $\theta = 10^\circ$, the asymmetry has been multiplied by 10.
b) The spin averaged cross-section in units of $\sigma(e^+e^- \rightarrow \mu^+\mu^-)$ at the same energy $\sqrt{s}$.

FIG. 1

FIG. 2

FIG. 3

FIG. 4

FIG. 5