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PROBING THE $WW\gamma$ VERTEX AT FUTURE HADRON COLLIDERS

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ABSTRACT

The structure of the $WW\gamma$ vertex can be measured in $W^\pm\gamma$ production at the LHC and the SSC. Deviations from the gauge theory structure can be parametrized in terms of 4 form factors which, if nonzero, would lead to additional events at large $W\gamma$ invariant masses. Signatures of these anomalous couplings for realistic experimental cuts are described and it is found that the two hadron colliders can improve the measurement of the $WW\gamma$ vertex by up to two orders of magnitude beyond what is expected via W^+W^- production in 190 GeV e^+e^- annihilation at LEP II.

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1. Introduction

During the past decade the standard electroweak theory (SM) has proven extremely successful in describing experimental results on electroweak interactions. In spite of this success one has to bear in mind that so far only the fermion-vector boson interactions have been scrutinized, while the scalar sector of the theory and the predictions of the SM for the vector boson self interactions have not yet been tested experimentally.

In this paper we concentrate on the vector boson self interactions and in particular on the $WW\gamma$ three vector boson vertex. Within the SM, at tree level, this vertex is completely fixed by the gauge theory structure of the model. Here we ask the question, how well will experiments at the Large Hadron Collider (LHC) (pp collisions at $\sqrt{s} = 17$ TeV) or at the Superconducting Supercollider (SSC) (pp collisions at $\sqrt{s} = 40$ TeV) be able to confirm the gauge theory nature of this vertex, or, alternatively, how sensitive would they be to deviations from the SM predictions.

The process that we investigate for this purpose, $W\gamma$ production from $q\bar{q}$ annihilation, has been studied by many authors in the past¹⁻⁵⁾, usually with emphasis on the anomalous magnetic moment of the W as a non gauge theory contribution. In this paper we go a step further and parametrize the most general $WW\gamma$ vertex which is accessible in the annihilation process $q\bar{q} \rightarrow W\gamma$ of effectively massless quarks. Four different anomalous couplings are allowed by electromagnetic gauge invariance and Lorentz invariance and we discuss their properties in Section 2. Apart from these anomalous contributions to the $WW\gamma$ vertex we assume the SM to be valid. In particular we assume the coupling of the electroweak bosons to quarks and leptons to be given by the SM.

Our analysis of the $W\gamma$ production process at hadron colliders is based on the calculation of helicity amplitudes for the complete process

$$q\bar{q} \rightarrow W^\pm \gamma, W^\pm \rightarrow l^\pm \nu.$$

In order to obtain realistic estimates for the sensitivity to anomalous couplings expected at the LHC or the SSC, we have imposed quite stringent cuts. These cuts are described in detail in Section 3, where a brief discussion of the background from W jet production, with the jet misidentified as a photon, can be found as well.

Already at present, severe constraints on anomalous $WW\gamma$ couplings exist, either derived from low energy experiments⁶⁻¹¹⁾, or from unitarity considerations^{7,12)}. In order to avoid confusion we shall first discuss the expected signals in $pp \rightarrow W^\pm \gamma X$, without taking

into account any low energy bounds on anomalous couplings (Section 4). However, we shall always respect tree level unitarity. Only afterwards shall we compare the sensitivity to anomalous couplings, expected at the SSC or the LHC, with the low energy bounds and with the sensitivity which is expected in W^+W^- production from e^+e^- annihilation at LEP II (Section 5).

2. $WW\gamma$ Couplings

W photon production from quark-antiquark annihilation depends on the structure of the $WW\gamma$ -vertex via the s -channel exchange of a W boson, as shown in Fig. 1. In $W\gamma$ production both the virtual W in the s -channel and the decaying onshell W couple to essentially massless fermions, which insures that effectively $\partial_\mu W^\mu = 0$. This together with gauge invariance of the onshell photon restricts the tensor structure of the WW photon vertex sufficiently to allow just four free parameters, which are conveniently described by the effective Lagrangian^{13,14)}

$$\begin{aligned} \mathcal{L}_{WW\gamma} = -ie \left\{ \left(W_{\mu\nu}^\dagger W^\mu A^\nu - W_\mu^\dagger A_\nu W^{\mu\nu} \right) + \kappa W_\mu^\dagger W_\nu F^{\mu\nu} + \frac{\lambda}{m_W^2} W_{\lambda\mu}^\dagger W_\nu^\mu A^{\nu\lambda} \right. \\ \left. + \tilde{\kappa} W_\mu^\dagger W_\nu \tilde{F}^{\mu\nu} + \frac{\tilde{\lambda}}{m_W^2} W_{\lambda\mu}^\dagger W_\nu^\mu \tilde{F}^{\nu\lambda} \right\}. \end{aligned} \quad (2.1)$$

Here A^μ and W^μ are the photon and W^- fields, respectively, $W_{\mu\nu} = \partial_\mu W_\nu - \partial_\nu W_\mu$, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, and $\tilde{F}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}F^{\rho\sigma}$. e is the charge of the proton.

The first term in Eq. (2.1) arises from minimal coupling of the photon to the $W^\pm$ fields and is completely fixed by the charge of the W boson for onshell photons. The κ and λ terms are related to the magnetic dipole moment μ_W and the electric quadrupole moment Q_W of the W^+

$$\mu_W = \frac{e}{2m_W} (1 + \kappa + \lambda), \quad (2.2a)$$

$$Q_W = -\frac{e}{m_W^2} (\kappa - \lambda). \quad (2.2b)$$

Within the standard model (SM), at tree level,

$$\kappa = 1, \quad (2.3a)$$

$$\lambda = 0. \quad (2.3b)$$

While κ and λ do not violate any discrete symmetries, the $\tilde{\kappa}$ and $\tilde{\lambda}$ terms are parity- and CP-odd. They are related to the electric dipole momenta d_W and the magnetic quadrupole moment $\tilde{Q}_W$ of the W^+ by

$$d_W = \frac{e}{2m_W}(\tilde{\kappa} + \tilde{\lambda}), \quad (2.4a)$$

$$\tilde{Q}_W = -\frac{e}{m_W^2}(\tilde{\kappa} - \tilde{\lambda}). \quad (2.4b)$$

Within the SM both vanish at tree level.

It is well known that tree level unitarity, *e.g.* for the process $e^+e^- \rightarrow W^+W^-$, uniquely restricts the $WW\gamma$ couplings to their (SM) gauge theory values at asymptotically high energies.¹⁵⁾ This implies that any deviation $a = \kappa - 1$, λ , $\tilde{\kappa}$, or $\tilde{\lambda}$ from the SM expectation has to be described by a form factor $a(\hat{s}, q_W^2, q_\gamma^2)$ which vanishes when either the $W\gamma$ invariant mass $\sqrt{\hat{s}}$ or the square of the four momentum of the final state W or photon (q_W^2 or q_γ^2) become large. In $W\gamma$ production $q_\gamma^2 = 0$ and $q_W^2 \approx m_W^2$, even when the finite W -width is taken into account. However, large values of $\hat{s}$ will be probed at future hadron colliders like the LHC or the SSC and the $\hat{s}$ dependence of the anomalous couplings has to be included in order to avoid unphysical results, which would violate unitarity. Consequently, throughout this paper, we use form factors

$$a(\hat{s}, q_W^2 = m_W^2, q_\gamma^2 = 0) = \frac{a_0}{(1 + \frac{\hat{s}}{\Lambda^2})^n}, \quad (2.5)$$

usually with the choice $n = 2$. We will comment later to what extent our results depend on this choice. Since terms proportional to λ and $\tilde{\lambda}$ (κ and $\tilde{\kappa}$) in the $W\gamma$ production amplitudes grow like $\hat{s}$ ($\sqrt{\hat{s}}$) in the high energy limit (see Eq. (2.7) below), $n > 1$ for $a = \lambda$ or $\tilde{\lambda}$ and $n > \frac{1}{2}$ for $a = \kappa$ or $\tilde{\kappa}$ would already guarantee unitarity in $q\bar{q} \rightarrow W\gamma$ at asymptotic scattering energies.

In Eq. (2.5) Λ represents the scale at which new physics becomes important in the weak boson sector, *e.g.* due to a composite structure of the W boson. The exponent $n = 2$ is chosen in order to suppress $W\gamma$ production at energies $\sqrt{\hat{s}} \gg \Lambda \gg m_W$, where novel phenomena like multiple weak boson production are expected to dominate. In Ref. 12 quantitative unitarity bounds on the low energy values a_0 were derived, assuming a common cutoff scale Λ in W^+W^- , $W\gamma$, and WZ production via fermion antifermion annihilation. Repeating this analysis with the form factor specified in Eq. (2.5), the following limits are

obtained for $n = 2$ and $\Lambda \gg m_W$:

$$|\kappa_0 - 1| < 7.4 \text{ TeV}^2/\Lambda^2, \quad (2.6a)$$

$$|\tilde{\kappa}_0| < 35 \text{ TeV}/\Lambda, \quad (2.6b)$$

$$|\lambda_0|, |\tilde{\lambda}_0| < 4.0 \text{ TeV}^2/\Lambda^2. \quad (2.6c, d)$$

When these bounds are observed tree level unitarity is satisfied throughout the entire $\hat{s}$ range, provided one anomalous coupling is nonvanishing at a time. For the more likely case that several anomalous couplings contribute, the bounds differ slightly and can be deduced from the results of Ref. 12.

The effective Lagrangian (2.1) may now be used to obtain cross-section formulas for $W\gamma$ production from $q\bar{q}$ annihilation, including the effect of any possible deviation of the $WW\gamma$ vertex from the SM prediction. Using amplitude calculation techniques, complete formulas have been given for W^+W^- production in Ref. 14 and for $W^\pm Z$ production at hadron colliders in Ref. 16, including the decay of the final state vector bosons into fermion-antifermion pairs, with or without radiation of a photon/gluon, and for general anomalous couplings. The analogous $W\gamma$ production amplitudes can easily be obtained by adapting the expressions of Ref. 16. We have used these amplitudes, which include a finite width W -propagator, for our numerical simulations. The program has been extensively checked. In particular electromagnetic gauge invariance for the emitted photon, up to effects proportional to the finite W -width in the propagator, has been verified for all anomalous couplings, which requires the inclusion of final state bremsstrahlung.

For our later phenomenological discussion it is convenient, however, to have available a rough approximation for the amplitudes and to display the contributions of anomalous couplings to the $q\bar{q} \rightarrow W\gamma$ helicity amplitudes in the high energy limit ($\hat{s} \gg m_W^2$). Neglecting the masses of the incoming quarks, the helicity of the quark (antiquark) is fixed to be $-\frac{1}{2}$ ($+\frac{1}{2}$) by the $V - A$ structure of the $Wq\bar{q}$ coupling. This means that the anomalous contributions to the $W\gamma$ production amplitudes only depend on the W and photon helicities λ_W and λ_γ . Denoting these contributions by $\Delta\mathcal{M}_{\lambda_\gamma\lambda_W}$ one finds

$$\begin{aligned} \Delta\mathcal{M}_{\pm 0} &= \frac{e^2}{\sin\theta_W} \frac{\sqrt{\hat{s}}}{2m_W} \left[\kappa - 1 + \lambda \mp i(\tilde{\kappa} + \tilde{\lambda}) \right] \frac{1}{2}(1 \mp \cos\Theta) \\ \Delta\mathcal{M}_{\pm\pm} &= \frac{e^2}{\sin\theta_W} \frac{\hat{s}}{2m_W^2} \left[\lambda \mp i\tilde{\lambda} \right] \frac{1}{\sqrt{2}} \sin\Theta, \end{aligned} \quad (2.7)$$

where Θ denotes the scattering angle of the photon with respect to the quark direction. Formulae for non-negligible $m_W/\sqrt{\hat{s}}$ can be found in Ref. 12.

The four helicity combinations in Eq. (2.7) are the only ones which get affected by anomalous couplings. The helicity combinations $(\lambda_\gamma, \lambda_W) = (+-)$ and $(-+)$ have the photon and W spins aligned along the photon momentum direction and hence have angular momentum $J \geq 2$: they cannot be reached by s -channel exchange of a vector boson. The fact that only the above four helicity combinations of the $J = 1$ partial wave can be reached by s -channel W exchange explains why four free parameters suffice to parametrize the effects of the most general $WW\gamma$ vertex in $W\gamma$ production.

While the SM contribution to the $q\bar{q} \rightarrow W\gamma$ scattering amplitudes is bounded from above for fixed scattering angle Θ , the anomalous contributions (2.7) rise without limit as $\hat{s}$ increases, eventually violating unitarity. This is the reason, of course, why the anomalous couplings have to show a form factor behaviour at high energies. On the other hand, when the amplitudes (2.7) approach the unitarity limit, they will completely dominate the much smaller SM contribution to the cross-section. Hence, at large $W\gamma$ invariant masses, Eq. (2.7) suffices to understand differential distributions of the photon and the W decay products.

3. Signal and Background

The signal we are investigating consists of an isolated, high p_T photon and an onshell W which may decay either leptonically or hadronically. It is obvious, that the hadronic decay modes of the W will be swamped by the QCD prompt photon background,¹⁷⁾ very much like W -pair production with one W decaying hadronically is essentially unobservable due to the W jet jet background.¹⁸⁾ By considering leptonic decay modes of the W only, roughly 20% of all W decays are still observable (we neglect the τ decay mode of the W in the following). We include the leptonic branching fraction

$$B = \text{Br} (W \rightarrow e\nu, \mu\nu) = 0.186. \quad (3.1)$$

(corresponding to a top-quark mass of 50 GeV) in all subsequent figures.

The experimental signal thus consists of an isolated, high p_T , electromagnetic cluster and an isolated charged lepton in association with missing p_T . We expect isolation cuts to effectively remove any background from heavy flavor production, *e.g.* $gg \rightarrow t\bar{t}\gamma$ with one top decaying semileptonically. On the other hand W jet production with the jet misidentified as a photon may cause a much more serious background. Such misidentification may arise due to fluctuations in the calorimetry or because the jet hadronizes, with a leading π^0

carrying away most of its energy. At SSC energies the probability for either fluctuations was estimated to be below 10^{-4} for jets of $p_T > 200$ GeV and 2.5×10^{-4} for $p_T > 100$ GeV from calorimetric information alone.¹⁹⁾ Since the thus misidentified photon will generally have a smaller p_T than its parent jet, and since no tracking information was included in the analysis of Ref. 19, a misidentification probability below 10^{-4} for $p_{T\gamma} > 100$ GeV appears to be a reasonable estimate. This value should also be compared with the misidentification probability of 4.6×10^{-4} measured by UA2 for $p_T > 50$ GeV jets at the CERN collider.²⁰⁾

In Fig. 2 we compare the p_T distribution of the signal ($pp \rightarrow W^+\gamma X, W^-\gamma X$) with the W jet background at $\sqrt{s} = 40$ TeV in the SM. Rapidity cuts $|y| < 2.5$ are imposed on W , photon, and jets in order to roughly simulate detector response. Since, in the case of misidentification, a jet of given p_T will always be viewed as a photon of lower p_T , a 10^{-4} misidentification probability will produce an even lower background than suggested by the figure. We conclude that a misidentification probability below 10^{-4} is sufficient to remove the W jet background. [The SM parameters used in Fig. 2 and all subsequent figures are $\alpha_{QED} = \alpha(m_W) = 1/128$, $m_Z = 93$ GeV, $m_W = 82$ GeV, and $\cos \theta_W = m_W/m_Z$. We use the distribution functions of Duke and Owens, set I.²¹⁾]

The misidentification problem will become more severe for small $p_{T\gamma}$ and it is not clear at present, down to which value the background rejection will be sufficient to obtain a clean signal. Hence we have imposed an, admittedly conservative, cut $p_{T\gamma} > 100$ GeV in all our simulations. In addition we impose a rapidity cut $|y_\gamma| < 2.5$, which, however, is only relevant for very energetic photons. The p_T cut removes both the collinear and the infrared singularity of the $q\bar{q} \rightarrow W\gamma$ subprocess. In order to simulate the performance of realistic detectors, for the decay products of the W we require $|y_{e,\mu}| < 3$, $p_{T_{e,\mu}} > 30$ GeV, and $p_{T,\text{miss}} > 30$ GeV. For W identification we further require the transverse mass m_{trans} reconstructed from the charged lepton and $p_{T,\text{miss}}$ to essentially agree with the W mass, $|m_{\text{trans}} - m_W| < 20$ GeV. This transverse mass cut together with the $p_{T\gamma}$ cut enforces a large angular separation of the photon and the charged lepton from W decay and effectively eliminates all final state bremsstrahlung events.

The effects of the cuts at the LHC and the SSC are summarized in Table 1. The most severe cut obviously is $p_{T\gamma} > 100$ GeV. In the next section we will discuss to which extent it limits the measurement of the $WW\gamma$ couplings. A few effects of the other cuts are noteworthy. The $p_{T,\text{miss}}$ cut of 30 GeV eliminates many more events than the equivalent cut on the charged lepton. This happens because the dominant $W^\pm$ helicity is $\lambda_W = \pm 1$,³⁾ implying that the charged lepton will tend to be emitted in the direction of the parent

W (see Fig. 8 for an illustration of the size of this effect), thus picking up most of its momentum. Hence, comparatively little of the W p_T is left for the neutrino.

The suppression of W^- versus W^+ events is largely due to the finite rapidity window for the charged leptons. According to the previous comment the W and charged lepton rapidities are strongly correlated and a cut on the W rapidity would have a very similar effect. It can be understood by noting that due to the different up- and down-quark charges the photon tends to be emitted in the direction of the original (anti)up-quark. Since no valence antiquarks are present in pp collisions, this implies that a W^+ tends to be emitted in the direction of the original low- x anti-down while a W^- moves in the direction of a larger- x valence down-quark. As a result the W^- receives a larger average lab frame rapidity, which makes it more likely to be lost in the beam pipe.

At large $W\gamma$ invariant masses $W^-\gamma$ production is further suppressed compared to $W^+\gamma$ production since the valence distribution for down-quarks decreases more rapidly with Feynman x than the up-quark valence distribution. Hence the luminosity function for $W^-\gamma$ production decreases more rapidly with increasing $W\gamma$ invariant mass. This has some effect on the observability of anomalous couplings in W^+ versus W^- production, as is obvious from the enhancement of anomalies by large values of $\sqrt{s} = m_{W\gamma}$ in Eq. (2.7).

It is clear that the cuts that we impose are somewhat arbitrary without simulating a definite detector, an enterprise that is beyond the purpose of this paper. Within our cuts and with an expected integrated luminosity of 10^4pb^{-1} after one or two years of running, the LHC (SSC) will provide about 500 (1200) clean $W^\pm\gamma$ events in a $p_{T\gamma}$ range which is particularly sensitive to anomalies in the $WW\gamma$ vertex, as we shall now see.

4. Signals for Anomalous Couplings

The prime advantage of high energy pp colliders is the very high center of mass energy $\sqrt{s}$ that is available in parton collisions. In the $q\bar{q} \rightarrow W\gamma$ subprocess, the effects of anomalies in the $WW\gamma$ vertex are greatly enhanced at large energies, due to the $\hat{\gamma} = \sqrt{s}/2m_W$ factors in the anomalous contributions to the amplitudes (2.7). Compared to $W\gamma$ production at the Tevatron^{4,5)} or W^+W^- production at LEP^{14,22)} this enhancement factor will particularly favor the observability of anomalous values of λ or $\tilde{\lambda}$, which are enhanced by $\hat{\gamma}^2$ in the amplitude whereas terms containing κ and $\tilde{\kappa}$ grow only linearly with $\hat{\gamma}$. This enhancement is limited by partial wave unitarity, of course.^{12,16)} As a rule of thumb, anomalous $WW\gamma$ photon couplings cannot enhance the $q\bar{q} \rightarrow W\gamma$ cross-section by more than a factor 10^3 ($\approx \sin^2\theta_W/\alpha^2$) over the SM cross-section at a fixed center of

mass energy $\sqrt{s}$, after cuts have been imposed to eliminate partial waves with high angular momentum, *i.e.* photons collinear to the initial state fermions. Unitarity of the $q\bar{q} \rightarrow W\gamma$ scattering amplitude will be insured in the following by using the form factors (2.5) with $n = 2$ for anomalous couplings.

A typical signal for anomalous couplings will be a broad increase in the W -photon invariant mass spectrum at large values of $m_{W\gamma} = \sqrt{s}$, possibly cut off above $m_{W\gamma} = \Lambda$, the scale of the form factor. In models of composite weak bosons²³⁾ one would expect additional resonance structure around $m_{W\gamma} \approx \Lambda$ due to the production and subsequent radiative decay of excited W 's.²⁴⁾ In the following we shall not consider such resonance structure; however, we shall allow for its existence above the form factor enhancement that we study by choosing the low energy values of the anomalous couplings well below their unitarity limits (2.6).

The resulting effect on $B \cdot d\sigma/dm_{W\gamma}$ is shown in Fig. 3 for a representative choice of anomalous couplings: $\kappa_0 = 1.5$, $\lambda_0 = 0.1$, $\bar{\kappa}_0 = 0.8$, and $\bar{\lambda}_0 = 0.05$, respectively, in the form factor (2.5) with $\Lambda = 2$ TeV and $n = 2$. Only one $WW\gamma$ coupling at a time is chosen different from the SM prediction. For comparison the SM curve is included as a solid line. The curves are shown both for the LHC (Fig. 3a) and the SSC (Fig. 3b).

At hadron colliders the $W\gamma$ invariant mass cannot unambiguously be determined, of course, due to the nonobservation of the neutrino arising from W -decay. Identifying the transverse momentum of the neutrino with the missing transverse momentum of a given $W\gamma$ event, the unobservable longitudinal neutrino momentum can be reconstructed, albeit with a twofold ambiguity, by imposing the constraint that the neutrino and charged lepton four-momenta combine to form the W rest mass.^{4,25)} What is shown in Fig. 3 is actually the "reconstructed" $W\gamma$ invariant mass spectrum obtained by using the two solutions for the neutrino longitudinal momentum for each event. Due to the transverse mass cut described in the last section, the reconstruction of $m_{W\gamma}$ is excellent and the "true" and "reconstructed" $W\gamma$ invariant mass distributions are almost identical.

In Fig. 3 the anomalous coupling curves are clearly distinguishable from the SM prediction, when remembering the expected luminosity of $10^4 \text{ pb}^{-1} \text{ year}^{-1}$ at both the SSC and the LHC: a differential cross-section of 10^{-7} pb/GeV corresponds to one event per year per TeV interval.

A pronounced feature of $W\gamma$ production in $q\bar{q}$ annihilation is the SM prediction of radiation zeros in all contributing helicity amplitudes at one value of the photon scattering angle and hence in the differential cross-section.²⁾ For $u\bar{d} \rightarrow W^+\gamma$ this radiation zero

occurs at $\cos \theta^* = -\frac{1}{3}$, where θ^* is the scattering angle of the photon relative to the u -quark direction, measured in the $W\gamma$ rest frame. Since the $W\gamma$ rest frame can be reconstructed, θ^* could be determined as well, if one knew which of the two incoming protons to associate with the u - and $\bar{d}$ -quarks, respectively. In practice, hence, only $|\cos \theta^*|$ is measurable and the radiation zero is considerably washed out.

In order to eliminate the strong peaking of the differential cross-section at $\cos \theta^* = \pm 1$ which arises from the collinear singularity, we study the rapidity distribution $d\sigma/d|y^*|$ of the photon instead, with

$$y^* = \frac{1}{2} \ln \frac{1 + \cos \theta^*}{1 - \cos \theta^*}. \quad (4.1)$$

The resulting $B \cdot d\sigma(pp \rightarrow W^+\gamma X)/d|y^*|$ distribution is shown in Fig. 4 for $m_{W\gamma} > 200$ GeV, but without any of the other cuts discussed in Section 3 (dashed line), and for our standard cuts (solid line). The dip at $|y^*| = 0$, which is due to the radiation zeros, is quite pronounced and remains even when taking into account the finite acceptance of detectors, as simulated by our cuts.

In the presence of any anomalous contribution to the $WW\gamma$ vertex the radiation zero will be eliminated. This is obvious from Eq. (2.7): none of the anomalous contributions to the scattering amplitudes vanishes at $\cos \theta^* = -\frac{1}{3}$. As a result the dip in the rapidity spectrum of the photon at $|y^*| = 0$ will be filled at least partially. This effect is shown in Fig. 5 for the same anomalous couplings and form factors as in Fig. 3, for both the LHC and the SSC.

Figure 5 clearly demonstrates that anomalous couplings mainly affect the region of small center of mass rapidities. This is due to the fact that anomalous couplings only contribute via the s -channel W -exchange graph of Fig. 1, and hence only to the $J = 1$ partial wave, when fermion masses are neglected. The anomalous contributions are, therefore, almost isotropic in the center of mass frame, while the u - and t -channel graphs of Fig. 1 result in the strong enhancement of the high rapidity region in Fig. 4. As a very fortunate result of this behaviour the finite acceptance cuts will largely eliminate the well known fermion exchange contributions to the cross-section and reduce possible signals of new physics, *i.e.* the effects of anomalous $WW\gamma$ interactions, by a much smaller amount.

The population of the small rapidity region, induced by anomalous couplings, considerably increases the average photon transverse momentum of events produced at a fixed value of the W photon invariant mass, as compared to the SM. The p_T distribution of the photon, $d\sigma/dp_{T\gamma}$, should hence be particularly sensitive to anomalous couplings. This

fact is clearly visible in Fig. 6 where the $p_{T\gamma}$ spectrum is plotted for the SM and for $\kappa = 1.5$, $\lambda = 0.1$, $\tilde{\kappa} = 0.8$, and $\tilde{\lambda} = 0.05$ and form factors as in Fig. 3. Experiments at both the SSC and the LHC would easily see the effect.

It is obvious from Fig. 6 that the cut $p_{T\gamma} > 100$ GeV has no appreciable effect on the observability of the small (and hence most interesting) anomalous couplings $\lambda = 0.1$ and $\tilde{\lambda} = 0.05$. Both the $1 + \cos^2\theta^*$ and $1 - \cos^2\theta^*$ distributions implied by the anomalous contributions to the production amplitudes (2.7) result in average $p_{T\gamma}$ values close to the kinematic limit of $\approx \sqrt{\hat{s}}/2$. As a result the cut $p_{T\gamma} > 100$ GeV only influences events and hence the measurement of form factors in the region $\sqrt{\hat{s}} = m_{W\gamma} \lesssim 300$ GeV. Measurements in this region are important only for small cutoff scales Λ , which damp anomalous coupling effects at larger $m_{W\gamma}$ values, and for anomalous couplings of order 0.3 or larger, in order to allow for sizable deviations from the SM predictions, even though the enhancement due to large values of $\sqrt{\hat{s}}/2m_W$ is not yet operative. If the scale of new physics Λ is small many other exciting effects, like multiple production of electroweak bosons, should produce striking signatures at both the LHC and the SSC, while large anomalous couplings will have been observed and studied earlier in W^+W^- production at LEP II. In view of this the transverse momentum cut on the photon appears as merely useful to eliminate unwanted SM events which can be described by the u - and t -channel fermion exchange graphs of Fig. 1.

As we have demonstrated so far, the transverse momentum and rapidity distributions of the produced photon, and the $W\gamma$ invariant mass spectrum are very sensitive indicators of anomalous couplings. We now want to quantify this statement, namely we shall derive the minimal values for κ , λ , $\tilde{\kappa}$, or $\tilde{\lambda}$ which would give rise to detectable effects at either the LHC or the SSC. Since the anomalous couplings are actually form factors, the answer depends to some extent on their functional behaviour. As before we choose a dipole form factor ($n = 2$ in Eq. (2.5)), and we explore the form factor dependence by varying the cutoff scale Λ . We use the following sensitivity criterion: an anomalous coupling is called detectable if it produces a deviation from the SM at the 99.99% confidence level in either the rapidity, the $m_{W\gamma}$ or the p_T spectrum of the photon calculated within our standard cuts. The confidence levels are calculated by splitting the $m_{W\gamma}$ and p_T spectra into 6 or 7 and the y^* distribution into 12 bins with, typically, more than 8 events. In each bin the Poisson statistics is approximated by a Gaussian distribution. In order to achieve a sizable counting rate in each bin, within the SM, all events with $p_T > 300$ GeV (400 GeV), or $m_{W\gamma} > 1$ TeV (1.7 TeV), are collected into a single bin at the LHC (SSC). This procedure guarantees that in our calculation a high confidence level cannot arise from a single event at

very high p_T , where the SM predicts, *e.g.* less than 10^{-2} events. The sensitivity estimates that we derive will thus be conservative. Furthermore we allow for a factor 2 normalization uncertainty of the SM prediction.

The resulting minimal anomalous couplings, which would give rise to a 99.99% CL effect are given in Tables 2a and 2b for the LHC and SSC for a range of cutoff scales Λ between 0.5 and 10 TeV. It turns out that in most cases the p_T spectrum of the photon is the most sensitive indicator of anomalous couplings. For $\Lambda \leq 1$ TeV, the y^* distribution is the most sensitive quantity, in particular for κ and $\tilde{\kappa}$. The $W\gamma$ invariant mass spectrum gives the highest confidence level for λ and $\tilde{\lambda}$ with $\Lambda \geq 10$ TeV. The values presented in Table 2 directly reflect the different powers of $\hat{\gamma} = \sqrt{\hat{s}}/2m_W$ multiplying the various anomalous contributions to the amplitudes (2.7). Since at high energies ($\hat{s} \gg m_W^2$) terms proportional to λ and $\tilde{\lambda}$ grow much faster than κ or $\tilde{\kappa}$ terms, the sensitivity limits achievable for λ and $\tilde{\lambda}$ are far better than the ones for κ and $\tilde{\kappa}$.

Table 2 has been derived using the $W^+\gamma$ signal only. The corresponding sensitivity bounds derived from the $W^-\gamma$ signal are up to 20% (40%) weaker for κ and $\tilde{\kappa}$ (λ and $\tilde{\lambda}$) because of the smaller counting rate in this channel. Since at pp colliders the $\bar{u}d$ luminosity decreases more rapidly with $m_{W\gamma}$ than the $\bar{d}u$ luminosity, the effect is particularly large at $\Lambda = 5$ and 10 TeV and for λ and $\tilde{\lambda}$.

Effectively the sensitivity to anomalous couplings in $pp \rightarrow W^\pm\gamma X$ stems from regions of phase space where the anomalous contributions to the cross-section are considerably larger than the SM expectation. As a result interference effects between the SM amplitude and the anomalous amplitude contributions (2.7) play a minor role, and an excess in counting rate, beyond the SM expectation, essentially scales like the sum of squares of anomalous couplings. This scaling can easily be observed in Fig. 6, for example. Over large regions of $p_{T\gamma}$ the cross-section for $\lambda = 0.1$ is almost precisely a factor 4 larger than the corresponding result for $\tilde{\lambda} = 0.05$. A similar statement holds for the $m_{W\gamma}$ spectrum of Fig. 3. This fact has a number of important consequences:

- i) The sensitivity bounds of Table 2 depend only marginally on the phases of the various anomalous couplings. Hence the limits are rather insensitive to the sign of the anomalous coupling. Furthermore, the limits for λ and $\tilde{\lambda}$ (κ and $\tilde{\kappa}$) are very similar because terms in the helicity amplitudes containing λ and $\tilde{\lambda}$ (κ and $\tilde{\kappa}$) differ, for $\hat{s} \gg m_W^2$, only by a factor $\pm i$. The marginal dependence of the differential cross-sections on the phases implies of course that it will also be difficult to measure phases of anomalous couplings, once they are found.

- ii) Contributions from different anomalous couplings cannot cancel at a significant level. Hence the sensitivities of Table 2, though derived by assuming only one anomalous coupling at a time to be different from the SM, represent model independent upper bounds on individual anomalous couplings that can be set by experiments.
- iii) Increasing statistics by a factor 2 will reduce the statistical error by a factor $\sqrt{2}$ and hence improve the sensitivity to any of the anomalous couplings by, roughly, a factor $\sqrt[4]{2}$, i.e. by less than 20%. An even smaller improvement is achieved by combining the $W^+\gamma$ and $W^-\gamma$ spectra. Incidentally, this is comparable to the estimated error of the bounds presented in Table 2.

The dependence of the sensitivity bounds in Table 2 on the cutoff scale Λ in the form factors has a simple explanation. It is illustrated in Fig. 7, where the $W^+\gamma$ invariant mass spectrum expected at the SSC is plotted for the SM (solid line) and for $\lambda_0 = 0.04$ with three choices for the form factor: $\Lambda = 1$ TeV and $n = 1$ (dashed line) and $n = 2$ (dotted line), and $\Lambda = 10$ TeV, $n = 2$ (dash-dotted line), which is at the limit of the unitarity bound (2.6). Obviously the improvement in sensitivity with increasing Λ is due to the additional events at large W photon invariant masses which were cut off by the form factor when the scale Λ had a smaller value.

Similarly the less drastic cutoff for $n = 1$ instead of $n = 2$ in the form factor allows for additional high $m_{W\gamma}$ events and therefore leads to an increased sensitivity to the low energy value λ_0 . For $\Lambda = 1$ TeV, for example, the SSC is sensitive to $\lambda_0 = 0.016$ (at the 99.99% CL) if $n = 1$ and to $\lambda_0 = 0.039$ if $n = 2$, which can be understood by noting that these two values of the low energy couplings give roughly the same value for the form factor in the critical energy region around $\sqrt{\hat{s}} = \Lambda$. It should be stressed that the difference between the $n = 1$ and $n = 2$ case becomes less pronounced for larger values of Λ , because the decrease in parton luminosity in the $\sqrt{\hat{s}} = \Lambda$ region shifts importance to the low energy region where differences in form factors are small. The sensitivities in Table 2 can thus be taken as representative for a much wider class of form factors.

From Eq. (2.6) it is obvious that unitarity forces the low energy values of the form factors to quickly decrease with increasing Λ , and, finally, to become so small that no deviation from the SM will be observable at future hadron colliders. The maximal scale one can probe at the LHC (SSC) with κ turns out to be $\Lambda \approx 6.2$ TeV (7.5 TeV) for a dipole form factor. The sensitivity limits given in Table 2 for $\Delta\kappa_0$ at $\Lambda = 10$ TeV, therefore, have no practical meaning and have merely been included for completeness. The corresponding maximal observable scale for λ and $\tilde{\lambda}$ is $\Lambda \approx 21$ TeV (32 TeV) whereas $\tilde{\kappa}$ formally would

allow to test Λ values in the 10^2 TeV region.* However, as we shall see in the next section, existing low energy bounds do not allow $\tilde{\kappa}$ to produce any visible effects on $W^\pm\gamma$ production at either the SSC or the LHC.

On the other hand, it is well possible that anomalous values of κ , λ , and/or $\tilde{\lambda}$ will produce clearly visible deviations from the SM. Thus the question arises to what extent experiments can distinguish these couplings. Inspection of the helicity amplitudes (2.7) shows that λ and $\tilde{\lambda}$ lead to an excess production of transverse W 's while a deviation of κ from its SM value 1 will enhance longitudinal W production.³⁾ The effect of these different W polarizations shows clearly in the angular distribution of the charged leptons arising from W decay. Figure 8 shows $B \cdot d\sigma(pp \rightarrow W^+\gamma X)/d\cos\bar{\theta}$ for the SM and the anomalous couplings and form factors as in Fig. 3. Here $\bar{\theta}$ is the e^+/μ^+ polar angle in the W^+ rest frame with respect to the W^+ direction in the $W\gamma$ rest frame. The cuts imposed on $p_{T,e,\mu}$ and $p_{T,\text{miss}}$ (see section 3) force the cross-section to vanish at $\cos\bar{\theta} = \pm 1$. Due to the $V-A$ structure of the $W e \nu$ coupling, transversely polarized W^+ 's of helicity $\lambda_W = \pm 1$ produce a $(1 \pm \cos\bar{\theta})^2$ distribution of the charged decay lepton, while longitudinal W 's result in a $\sin^2\bar{\theta}$ distribution. The $\cos\bar{\theta}$ distributions of Fig. 8 exhibit the pronounced dominance of left-polarized W 's in the standard model, of longitudinal W 's for large anomalous values of κ ,³⁾ and of equal numbers of left- and right-handed W 's for anomalous values of λ and $\tilde{\lambda}$. The W decay distribution will thus allow a separation of κ from λ and $\tilde{\lambda}$. A separation of the latter two requires the measurement of interference terms between anomalous and SM contributions to individual scattering amplitudes or between helicity amplitudes corresponding to the same photon but different W polarizations. Since these interference terms can easily change sign as a function of center of mass energy and/or scattering angle, their measurement will most likely require the investigation of correlations of observables, which may well prove prohibitive in view of the limited statistics of clean $W\gamma$ events.

Finally, we would like to briefly discuss the $p_{T,\text{miss}}$ spectrum. From Fig. 8 it is obvious that, in the SM, the charged lepton tends to be emitted in the direction of the parent W . It will, therefore, pick up most of the W momentum, and only relatively little p_T is left for the neutrino. As a result, the SM $p_{T,\text{miss}}$ distribution drops much faster than the p_T spectrum of the charged lepton. This is not the case if large anomalous couplings are present. Since the contributions of anomalous couplings to the $\cos\bar{\theta}$ distribution are

* For $n = 1$ instead of $n = 2$ in the form factor the unitarity limits for $\Delta\kappa_0$, λ_0 and $\tilde{\lambda}_0$ are directly given by the bounds listed in Ref. 12 and the maximal scale one can probe with κ , λ and $\tilde{\lambda}$ is reduced by about a factor 2. The range which could be explored with $\tilde{\kappa}$ is only little changed.

all symmetric with respect to $\cos \bar{\theta}$, the charged lepton and the neutrino will have rather similar p_T distributions. This means that the effect of the anomalous couplings will be more pronounced in the $p_{T,\text{miss}}$ spectrum than in the $p_{T_{e,\mu}}$ distribution. Although the missing transverse momentum spectrum is very sensitive to deviations from the SM, we have not included this quantity to derive the sensitivity bounds in Table 2 for the following reason. For anomalous couplings close to the sensitivity limits the $p_{T,\text{miss}}$ spectrum still drops much faster than the $p_{T_{e,\mu}}$ distribution, and a very good energy resolution will be required for the calorimeters, in order to improve the limits of Table 2. Leptonic τ decays may degrade the $p_{T,\text{miss}}$ spectrum even further. Even if one takes into account the $p_{T,\text{miss}}$ distribution in addition to the $m_{W\gamma}$, $p_{T\gamma}$, and y^* spectra, the sensitivity bounds in Table 2 do not improve by more than 10%, as we have verified.

5. Discussion and Conclusions

We have described, in the last section, the signatures which anomalies in the $WW\gamma$ vertex will produce in $W\gamma$ production at future hadron colliders, and in particular we have determined how large deviations from the SM have to be in order to yield visible effects in experiments at the LHC or the SSC. Naturally the question arises how the sensitivity of these two hadron colliders compares with existing low energy limits on anomalous couplings and with the sensitivity to non-gauge theory terms in the three boson vertex reachable at LEP by investigating the process $e^+e^- \rightarrow W^+W^-$. It must be realized, of course, that these three approaches to determining anomalous couplings measure the form factors corresponding to the couplings in quite different regions of momentum space: timelike momenta in the vector boson production processes and spacelike momenta in the low energy bounds, which are derived by performing loop calculations involving the $WW\gamma$ vertex. In this sense the three approaches are complementary. Since we have parametrized the form factors κ , λ , *etc.* in terms of the small momentum transfer values κ_0 , λ_0 , *etc.* which necessarily are the same in all three approaches, low energy bounds and bounds from $e^+e^- \rightarrow W^+W^-$ must nevertheless be incorporated by the time the hadron collider experiments are performed. Of course these bounds will not rule out the possible existence of resonances in the $W\gamma$ channel, which also can be parametrized via these form factors.

Low energy bounds on κ and λ have been derived by calculating loop contributions to $(g-2)_\mu$,⁶⁻⁸⁾ the photon propagator as measured at PETRA,⁸⁾ and the W/Z mass ratio.^{7,8)} The W/Z mass ratio produces the most stringent, but also the most controversial bounds.^{7,8)} Since any non-gauge model of the W and the Z has to explain their small mass

(as compared to the compositeness scale Λ) and their mass ratio, probably by some as yet unknown symmetry between the various $WW\gamma$ and WWZ couplings, bounds on individual anomalous couplings derived from the W and Z masses should not be used to predict null results in high energy scattering experiments. The most stringent bound is then obtained from $(g-2)_\mu$, namely⁶⁻⁸⁾

$$\left| (\kappa - 1) \ln \frac{\Lambda^2}{m_W^2} + \lambda \right| < 4.4 \quad (5.1)$$

which, for cutoff scales larger than a few TeV, gets surpassed by the unitarity bounds (2.6). From Table 2 it is obvious that experiments at the LHC and the SSC will considerably improve the present low energy bounds on κ and λ . This statement is not true for the CP violating coupling $\tilde{\kappa}$, which produces a potentially large contribution to the electric dipole moment of the neutron. In Ref. 10 the present experimental limit on the neutron electric dipole moment was used to derive the bound

$$\left| \tilde{\kappa} \ln \frac{\Lambda^2}{m_W^2} \right| < 10^{-3}, \quad (5.2)$$

which clearly excludes any observable effect in $W\gamma$ production, as shown by a comparison with Table 2. Curiously no such bound exists for the other CP-violating coupling $\tilde{\lambda}$. Contributions of $\tilde{\lambda}$ to the neutron electric dipole moment are suppressed by at least another factor of $(m_n/m_W)^2$ (m_n being the neutron mass),¹¹⁾ as compared to the $\tilde{\kappa}$ bound (5.2), and hence experiments at the LHC and the SSC will explore a range of three orders of magnitude in $\tilde{\lambda}$.

Experiments at LEP, studying W^+W^- production in e^+e^- annihilation at 190 GeV, can measure κ , λ , $\tilde{\kappa}$, and $\tilde{\lambda}$ with an accuracy of $\pm 0.2 \dots \pm 0.7$, depending on the values of the interfering WWZ couplings.²²⁾ While it will be difficult for the future hadron colliders to improve the measurement of κ in W photon production, the determination of λ and $\tilde{\lambda}$ can be improved by up to two orders of magnitude.

The expected sensitivities to $WW\gamma$ anomalous couplings that we have calculated are somewhat better (by at most a factor 2), than the corresponding sensitivities to anomalous WWZ couplings in WZ production at the LHC or the SSC.¹⁶⁾ The larger sensitivity in $W\gamma$ production is partially due to the larger counting rate at high p_T (the limiting factor in WZ production is the small branching fraction of the Z into electron or muon pairs). Partially it arises because the SM prediction of radiation zeros produces additional structure in $W\gamma$ differential cross-sections which is not available in the WZ channel.

In spite of the high sensitivity to λ and $\tilde{\lambda}$, neither the LHC nor the SSC will be able to test the radiative corrections to these quantities in the SM, which are at most of $O(0.01)$. The reason for this is that, in the SM, the form factors vary on a scale given by the masses of particles in the loop diagrams, *e.g.* the weak boson masses or heavy fermion masses. These are, normally, at most in the few hundred GeV region, and according to Table 2 future pp colliders are not able to limit λ and $\tilde{\lambda}$ to better than $O(0.1)$ for scales in this range.

It is quite instructive to relate the accuracy of the various methods for determining non-gauge theory $WW\gamma$ interactions with the magnitude of the scale of new physics, Λ , which can be probed. Assuming that possible deviations are due to some novel strong interactions in the electroweak boson sector, the anomalous couplings may be expected close to their unitarity limits (2.6). Saturating these unitarity limits and assuming dipole form factors for all the anomalous couplings, the low energy bounds (5.1) probe scales up to $\Lambda \approx 2$ TeV at most, LEP is sensitive up to $\Lambda \approx 4$ -5 TeV while the SSC (LHC) probes scales up to 32 TeV (21 TeV) in $W\gamma$ production. Thus W photon production at both the SSC and the LHC provides a tool to investigate the properties of the weak boson sector well above the few TeV region where the production of new particles may yield a more direct signature.

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TABLE 1

Effects of increasingly stringent cuts on the expected number of events $\int \mathcal{L} dt B \cdot \sigma(pp \rightarrow W\gamma X)$ in the SM. In each line the indicated cut is imposed in addition to all the previous ones. An integrated luminosity of 10^4pb^{-1} is assumed.

$\sqrt{s}$ process	17 TeV		40 TeV	
	$pp \rightarrow W^+\gamma X$	$pp \rightarrow W^-\gamma X$	$pp \rightarrow W^+\gamma X$	$pp \rightarrow W^-\gamma X$
$p_{T\gamma} > 50 \text{ GeV}$	10200	7540	25100	19600
$p_{T\gamma} > 100 \text{ GeV}$	1600	1110	4410	3350
$ y_\gamma < 2.5$	1080	1010	2620	2750
$ p_{T_{e,\mu}} > 30 \text{ GeV}$	920	900	2250	2470
$ y_{e,\mu} < 3$	890	770	2000	1830
$ p_{T,\text{miss}} > 30 \text{ GeV}$	560	470	1260	1130
$ m_{\text{trans}} - m_W < 20 \text{ GeV}$	280	240	640	590

TABLE 2

Sensitivities achievable at the 99.99% CL for the four anomalous $WW\gamma$ couplings a) at the LHC and b) at the SSC. Dipole form factors with a cutoff scale Λ are assumed. (See text for details.) The bounds are derived from the $W^+\gamma$ channel only. The values for $\Delta\kappa_0$ at $\Lambda = 10$ TeV are larger than what is allowed by unitarity and, therefore, have been included in parentheses.

a) LHC bounds ($\sqrt{s} = 17$ TeV, $\int \mathcal{L} dt = 10^4 \text{pb}^{-1}$)

Λ	0.5 TeV	1 TeV	2 TeV	5 TeV	10 TeV
$\Delta\kappa_0$	± 0.58	± 0.34	± 0.27	± 0.20	(± 0.18)
λ_0	± 0.12	± 0.054	± 0.021	± 0.012	± 0.010
$\tilde{\kappa}_0$	± 0.54	± 0.32	± 0.26	± 0.20	± 0.18
$\tilde{\lambda}_0$	± 0.12	± 0.052	± 0.021	± 0.012	± 0.010

b) SSC bounds ($\sqrt{s} = 40$ TeV, $\int \mathcal{L} dt = 10^4 \text{pb}^{-1}$)

Λ	0.5 TeV	1 TeV	2 TeV	5 TeV	10 TeV
$\Delta\kappa_0$	± 0.46	± 0.26	± 0.20	± 0.14	(± 0.13)
λ_0	± 0.095	± 0.039	± 0.014	± 0.0066	± 0.0046
$\tilde{\kappa}_0$	± 0.43	± 0.26	± 0.20	± 0.13	± 0.12
$\tilde{\lambda}_0$	± 0.094	± 0.039	± 0.013	± 0.0065	± 0.0047

FIGURE CAPTIONS

1. Feynman graphs for the subprocess $q\bar{q} \rightarrow W\gamma$.
2. Photon p_T spectrum for the signal $pp \rightarrow W^\pm\gamma X$ (dashed line) and the background $pp \rightarrow W^\pm \text{ jet } X$ (solid line). The background is multiplied by a factor 10^{-4} to represent the misidentification probability of a jet as a photon. Rapidity cuts $|y| < 2.5$ are imposed on the W , the photon and the jet.
3. W photon invariant mass spectrum for the process $pp \rightarrow W^+\gamma X$ at (a) the LHC and (b) the SSC. Cuts are specified in the text. The curves are for the SM (solid line), $\kappa = 1.5$ (dashed line), $\lambda = 0.1$ (long-dashed line), $\tilde{\kappa} = 0.8$ (dot-dashed line) and $\tilde{\lambda} = 0.05$ (dotted line). Dipole form factors ($n = 2$ in Eq. (2.5)) with $\Lambda = 2$ TeV are chosen for all anomalous couplings.
4. Rapidity spectrum of the photon in the $W\gamma$ rest frame for the process $pp \rightarrow W^+\gamma X$ at $\sqrt{s} = 40$ TeV. The dashed line shows the spectrum in the narrow width approximation without any cuts but $m_{W\gamma} > 200$ GeV, multiplied by 0.1. The solid line is the portion surviving the cuts described in the text.
5. Rapidity spectrum of the photon in $pp \rightarrow W^+\gamma X$ in the $W\gamma$ rest frame at the (a) LHC and (b) SSC. Anomalous couplings, form factors and cuts are the same as in Fig. 3.
6. Transverse momentum spectrum of the photon in $pp \rightarrow W^+\gamma X$. Parameters and cuts are chosen as in Fig. 3.
7. $W^+\gamma$ invariant mass spectrum at the SSC for the SM (solid line) and $\lambda_0 = 0.04$ with three choices of form factors: $n = 1$, $\Lambda = 1$ TeV (dashed line); $n = 2$, $\Lambda = 1$ TeV (dash-dotted line) and $n = 2$, $\Lambda = 10$ TeV (dotted line) in the form factor of Eq. (2.5). Cuts are specified in section 3.
8. Polar angle distribution of the charged decay lepton from W^+ decay in $pp \rightarrow W^+\gamma X$ at the (a) LHC and (b) SSC. See text for definition of $\bar{\theta}$. Cuts and anomalous couplings are the same as in Fig. 3.

Fig. 1

Fig. 2

Fig. 3a

Fig. 3b

Fig. 4

Fig. 5a

Fig. 5b

Fig. 6a

Fig. 6b

Fig. 7

Fig. 8a

Fig. 8b