Direct photon production in $\gamma\gamma$ collisions

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The production of a hard photon balanced by a jet in the reaction $\gamma+\gamma\to\text{jet}+\gamma+X$ is shown to offer a good probe of the hadronic structure of the photon at e^+e^- colliders. The cross section for $p_T^\gamma > 1.5$ GeV ranges from ~ 0.3 pb at PEP to ~ 1 pb at TRISTAN to ~ 6 pb at LEP 200. For $\sqrt{s}\geq 150$ GeV, the study of this process will yield important information on the gluon density inside the photon, about which very little is known experimentally at present.

In recent papers [1,2] we studied the production of two high-p_T jets in almost-real $\gamma\gamma$ collisions [3]. The idea was to use these processes, which can be studied at the existing e^+e^- colliders TRISTAN [1] and LEP [2], to extract information on the quark and gluon content of the photon. At present, the former has been studied experimentally [4] via deep inelastic scattering (DIS), only in the region $x\geq 0.05$ and $Q^2\lesssim 100\ \text{GeV}^2$, while almost nothing is known about the latter. Information on the parton densities $\vec{q}^\gamma(x, Q^2)$, with $\vec{q}^\gamma\equiv(q_i^\gamma, G^\gamma)$, could then be used to sharpen predictions for photoproduction experiments at fixed targets [5,6] or at HERA [7] as well as for photon initiated cosmic air showers [8].

The main advantage of studying jet production in real $\gamma\gamma$ collisions rather than the "classic" DIS process [4] where at least one of the two photons is far off-shell lies in the fairly large cross sections, which are $O(\sim 100\ \text{pb})$ for TRISTAN energy ($\sqrt{s}=60$ GeV) and $p_T=3$ GeV. the disadvantage is that the only kinematic variables that are available are the transverse momenta and rapidities of the two "high-p_T" jets. Since the cross sections become uninterestingly small for $p_T(\text{jet})\geq 10$ GeV, hadronisation effects could be sizeable. Note that these events contain up to two spectator jets (see below) in addition to the two high-p_T jets, with nontrivial colour flow between spectator and hard jets. In addition, the direct process $\gamma\gamma\to q\bar{q}$ gives a large potential background to the processes that probe the hadronic structure of the photon.

Both these two problems are alleviated if we replace one of the two high-p_T jets by an isolated photon. In this letter we therefore study direct photon production in quasi real $\gamma\gamma$ collisions at e^+e^- colliders:

$$e^+e^-\to e^+e^-\gamma\,\text{jet}+X\,, \qquad (1)$$

where the hard jet balances the p_T of the outgoing photon. We are interested in no-tag events, where the outgoing e^+ and e^- are not detected. Note that in leading order this hard scattering necessarily involves the quark or gluon content of at least one photon. It should be noted here that inclusion of the process $\gamma\gamma\to q\bar{q}\gamma$ would be double counting. If we require the emitted photon to (almost) balance the transverse momentum of one of the jets, the exchanged quark has to be (almost) on-shell. The corresponding collinear divergence has to be absorbed in the distribution function $q_i^\gamma(x, Q^2)$. The process $\gamma\gamma\to q\bar{q}\gamma$ would therefore only contribute if we require *two* high-p_T jets. This leaves us with the contributions depicted in fig. 1: The "single resolved" process of fig. 1a, where one of the photons Compton-scatters off a quark within the other photon, and the two "double resolved" contributions of fig. 1b, where the hard processes are $q\bar{q}\to g\gamma$ and $g\bar{q}\to\gamma\bar{q}$, respectively. We note here that this last process provides an important tool [9] for constraining the gluon content of had-

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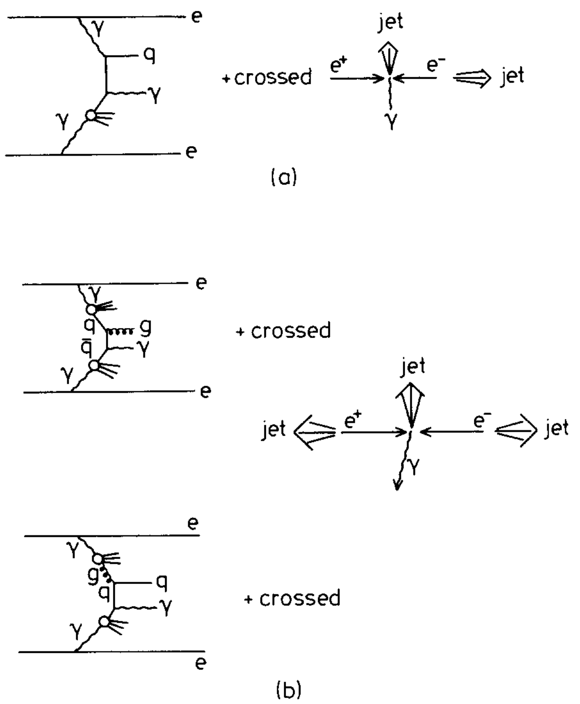

Fig. 1. The Feynman diagrams (left) and final state topologies (right) of the process of eq. (1). Single- and double-resolved contributions are shown in (a) and (b). Note that the spectator jet of the single-resolved process could also emerge in the direction of the e$^-$ beam.

rons via direct photon production in hadronic collisions. Fig. 1 also shows the final-state topology of the two classes of contributions. Each resolved photon produces a spectator jet, summarily denoted by "X" in eq. (1), which tends to go down the beam pipe.

The general expression for the cross section can be written as

$$\frac{\mathrm{d}\sigma}{\mathrm{d}p_T} = 2p_T \int\limits_{-y_{1,\max}}^{y_{1,\max}} \mathrm{d}y_1 \int\limits_{y_{2,\min}}^{y_{2,\max}} \mathrm{d}y_2 \times \int\limits_{x_{1,\min}}^{1} \frac{\mathrm{d}x_1}{x_1} f_{\gamma/e}\left(\frac{x_{1,\min}}{x_1}\right) \bar{q}^\gamma(x_1, Q^2) \times \int\limits_{x_{2,\min}}^{1} \frac{\mathrm{d}x_2}{x_2} f_{\gamma/e}\left(\frac{x_{2,\min}}{x_2}\right) \bar{q}^\gamma(x_2, Q^2) \frac{\mathrm{d}\hat{\sigma}}{\mathrm{d}\hat{t}}(\hat{s}, \hat{t}, \hat{u}), \tag{2}$$

where y_1 and y_2 are the rapidities of the two outgoing hard particles and x_1 and x_2 are the Bjorken-x variables for the partons inside the photons. The integration boundaries are given by

$$y_{1,\max} = \ln[1/x_T + \sqrt{(1/x_T^2 - 1)}], \tag{3a}$$

$$y_{2,\min} = -\ln[2/x_T - \exp(-y_1)], \tag{3b}$$

$$y_{2,\max} = \ln[2/x_T - \exp(y_1)], \tag{3c}$$

$$x_{1,\min} = \tfrac{1}{2}x_T[\exp(y_1) + \exp(y_2)], \tag{3d}$$

$$x_{2,\min} = \tfrac{1}{2}x_T[\exp(-y_1) + \exp(-y_2)], \tag{3e}$$

where $x_T = 2p_T/\sqrt{s}$. The Mandelstam variables for the hard process are given by

$$\hat{s} = s x_{1,\min} x_{2,\min}, \tag{4a}$$

$$\hat{t} = -\tfrac{1}{2}\hat{s}(1 \pm \sqrt{1 - 4p_T^2/\hat{s}}), \tag{4b}$$

with $\hat{u} = -\hat{s} - \hat{t}$. Note that both signs in eq. (4b) have to be included, since the photon and the jet are physically distinguishable. Alternatively one can fix the sign in eq. (4b) and add the $\hat{t} \leftrightarrow \hat{u}$ exchanged version to the square of the matrix element $\mathrm{d}\hat{\sigma}/\mathrm{d}\hat{t}$. For the single resolved process one has to either replace $\bar{q}^\gamma(x_1, Q^2)$ or $\bar{q}^\gamma(x_2, Q^2)$ with $\delta(1 - x_{1,2})$. Finally, for completeness we list the three relevant matrix elements [9]:

$$\frac{\mathrm{d}\hat{\sigma}(\gamma q \to \gamma q)}{\mathrm{d}\hat{t}} = \frac{2\pi\alpha^2}{\hat{s}^2} e_q^4 \left(-\frac{\hat{s}}{\hat{u}} - \frac{\hat{u}}{\hat{s}}\right), \tag{5a}$$

$$\frac{\mathrm{d}\hat{\sigma}(q\bar{q} \to \gamma g)}{\mathrm{d}\hat{t}} = \frac{8\pi\alpha\alpha_s}{9\hat{s}^2} e_q^2 \left(\frac{\hat{t}}{\hat{u}} + \frac{\hat{u}}{\hat{t}}\right), \tag{5b}$$

$$\frac{\mathrm{d}\hat{\sigma}(gq \to \gamma q)}{\mathrm{d}\hat{t}} = \frac{\pi\alpha\alpha_s}{3\hat{s}^2} e_q^2 \left(-\frac{\hat{s}}{\hat{u}} - \frac{\hat{u}}{\hat{s}}\right). \tag{5c}$$

Notice that, even though the single resolved matrix element (5a) is suppressed by a factor α/α_s compared to the processes (5b) and (5c), the two contributions are in fact of the same order in coupling constants, since the parton densities inside the photon are [10] of order α/α_s.

In eq. (2) we have used the Weizsäcker–Williams or the effective photon approximation [11], where the cross section for process (1) is calculated as a convolution of real γγ cross sections with the flux factors $f_{\gamma/e}$ [12]. We have used for these functions the form given in ref. [12], since this includes "finite terms" in addition to the leading logarithms, and has been shown [13] to describe both, total *and* differential cross sections for lepton pair production in $\gamma\gamma$ collisions to better than 5%.

At present only two parametrisations [6,14] of parton densities inside the photon exist. Since the parametrisation of ref. [14] is valid for $\Lambda_{\mathrm{QCD}}=400$ MeV, we will use this value throughout. As discussed in some detail in ref. [1], the "asymptotic", DO parametrisation of ref. [6] then has to be augmented by some "hadronic" contribution in order to describe existing DIS data [15] on the electromagnetic structure function F_2^γ of the photon; we add a contribution inspired by the vector meson dominance model (VDM), as discussed in ref. [1]. We use $Q^2=p_T^2$ both in the parton densities and in α_s. Notice that, unlike in hadronic collisions, the choice of a larger Q^2 scale will tend to *increase* the cross section, since $\vec{q}^\gamma(x, Q^2)$ has a component that increases logarithmically [10,14] with Q^2. E.g., for $\sqrt{s}=60$ GeV and $p_T=2$ GeV, the prediction of the DG parametrisation [14] for the single-resolved contribution increases by $\sim 25\%$ if we use $Q^2=\hat{s}$, while the double-resolved contribution remains almost unchanged. Finally, since the cross section is sizeable only for p_T not much larger than the charm mass, we use $N_f=3$ massless flavours everywhere; the charm contribution can be computed from $\gamma\gamma\to c\bar{c}\gamma$, but will in general have a different topology from the contribution of the lighter flavours, as there is no large (logarithmic) enhancement for collinear $c\bar{c}$ production.

In figs. 2a, 2b we present results for the transverse momentum spectrum of the produced photon at $\sqrt{s}=60$ GeV, for the two parametrisations discussed above. We have restricted ourselves to the region $p_T>1.5$ GeV since for even smaller values of the transverse momentum, QCD predictions clearly become unreliable. We see that the two parametrisations make very similar predictions for the double resolved contribution from $q\bar{q}$ initial states, as well as for the single resolved contribution. This is not surprising, since the VMD component in the DO+VMD parametrisation has been normalised [1] such that F_2^γ, which is directly proportional to the quark densities, is very similar for the two parametrisations for $Q^2=5$ GeV^2. Notice that the unphysical $x\to 0$ divergence of the "asymptotic" DO parametrisation increases the single resolved contribution only by 10% even at the lowest p_T of 1.5 GeV. The reason is that, although in principle values of Bjorken-x as low as

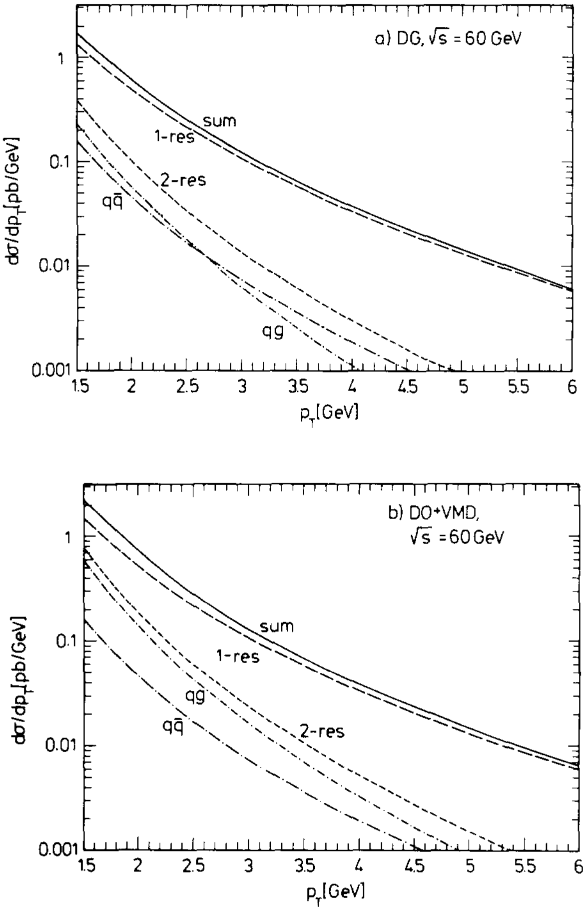

Fig. 2. The transverse momentum spectrum at $\sqrt{s}=60$ GeV, for the parametrisation of ref. [14] (a) and the sum of the "asymptotic" parametrisation of ref. [6] and a VMD inspired soft part (b), as described in the text. Results are for $\Lambda_{\mathrm{QCD}}=400$ MeV, $N_f=3$ flavours and $Q^2=p_T^2$, in the effective photon approximation of ref. [12]. The long-dash-dotted and short-dash-dotted curves represent the two double-resolved contributions with $q\bar{q}$ and $\bar{q}g$ initial states, see fig. 1. The short-dashed, long-dashed and solid curves show the total double- and single-resolved contributions and the sum of both, respectively. All calculations are in leading order, so that the transverse momenta of the high-p_T jet and the outgoing photon are equal and opposite.

3×10^{-3} can be probed at this value of p_T, the bulk of the cross section comes from much larger values of x, due to the softness of the photon density function $f_{\gamma/e}$.

It should also be noted that these contributions are quite insensitive to the flavour structure of the photon. Here we have assumed an SU(3) symmetric

VMD contribution. Had we assumed this "soft" contribution to q_i^γ to be proportional to $e^2_{q_i}$ (which is true for the "asymptotic" contribution at medium and large x), the single resolved "VMD" contribution would only have increased by a factor of $\frac{11}{9}$, once we require $F_2^{\gamma\mathrm{MD}}$ to be fixed by the data. This means that the total single-resolved contribution would have increased only by about 10%, which is hardly significant compared to the uncertainties due to the higher order QCD corrections, etc.

In sharp contrast, the predictions for the two parametrisations for the double-resolved contributions for the qg initial state differ by a factor 2.5 or more at all p_T values of interest. This is due to the large gluon component of the VMD part of the DO+VMD parametrisation; in the DG parametrisation, G^γ is essentially only created radiatively [14,16]. Since the qg initial state dominates the total double-resolved contribution at least for $p_\mathrm{T}^\gamma \leqslant 2.5$ GeV, the difference in G^γ leads to a fairly large ($\sim 100\%$) uncertainty in this quantity. This uncertainty becomes somewhat smaller at larger p_T^γ, since the qg initial state is relatively less important at higher transverse momentum due to the softness of $G^\gamma(x)$. We also see from fig. 2 that given an integrated luminosity of a few hundred pb^{-1}, at TRISTAN one can probe the single resolved process up to $p_\mathrm{T}^\gamma \simeq 5.5$–6 GeV. The double-resolved contribution is in principle detectable for $p_\mathrm{T}^\gamma \leqslant 3$ (3.5) GeV for the DG (DO+VMD) parametrisation.

In figs. 3a, 3b we show the results for the triple differential cross section $\mathrm{d}\sigma/\mathrm{d}p_\mathrm{T}\mathrm{d}y_\gamma\mathrm{d}y_\mathrm{jet}$ at $p_\mathrm{T}=2$ GeV, $y_\mathrm{jet}=0$ and $\sqrt{s}=60$ GeV. We again see that the two parametrisations make very similar predictions for the single-resolved and the $\mathrm{q\bar{q}}$ initiated double-resolved contributions, but differ quite strongly for the double-resolved qg-initiated contribution. We also see that this last contribution is somewhat more strongly peaked at the symmetric point $y_\gamma = y_\mathrm{jet}$ ($=0$ in our example), where the γ–jet invariant mass is minimal; this is again a reflection of the softness of $G^\gamma(x)$. However it is doubtful whether this rather small difference will be useful for disentangling the two double-resolved contributions from each other.

Potentially a more useful tool for this purpose could be the energy of the spectator jet in the hemisphere opposite to the outgoing photon. This energy is proportional to $(1-x_1)$, where x_1 is the Bjorken-x of the parton inside the initial-state photon that goes in the hemisphere opposite to the final-state photon. By requiring fairly large $|y_\gamma|$, say $|y_\gamma| \geqslant 1.5$, we can force the Bjorken-x (x_2) of the parton in the *other* initial-state photon to be large (see eqs. (3d), (3e)); the softness of G^γ then implies that this parton will almost always be a quark. In addition, it means that x_1 will be small, i.e. $(1-x_1)$ and hence the spectator jet energy will be large. Indeed we find that for $|y_\gamma| \geqslant 1.5$, the average energy of this spectator jet is about 10–12 GeV for the qg-initiated contribution, but only 5–6 GeV for the $\mathrm{q\bar{q}}$-initiated contribution. Unfortunately in both cases a good fraction of this energy will go down the beam pipe; the question whether this factor of two difference in the average energies is de-

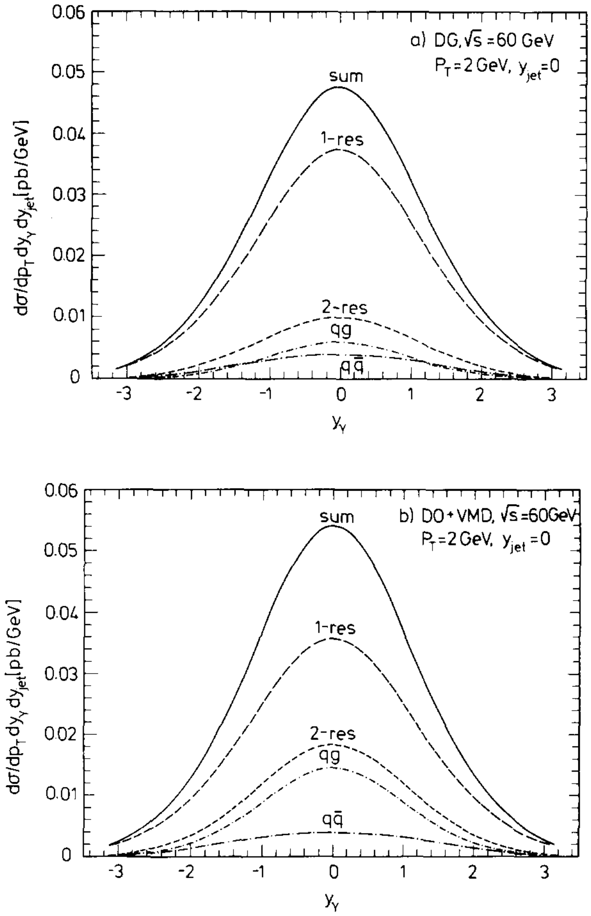

Fig. 3. The triple differential cross section $\mathrm{d}\sigma/\mathrm{d}p_\mathrm{T}\mathrm{d}y_\mathrm{jet}\mathrm{d}y_\gamma$ at $\sqrt{s}=60$ GeV, $p_\mathrm{T}=2$ GeV and $y_\mathrm{jet}=0$, for DG (a) and DO+VMD (b) parametrisations. Parameters and notation are as in fig. 2.

tectable therefore depends crucially on the angular coverage of the detector.

So far we have focussed on the centre-of-mass energies relevant for the TRISTAN collider. What are the prospects of studying this process at other e^+e^- colliders, past or future? To answer this question, we show in fig. 4 the cross section (2) integrated over the region $p_T > 1.5$ GeV as a function of $\sqrt{s}$. We see that about 100 events of this type might have been detected by each PEP and PETRA experiment; recall that the historic first measurement of F_2^γ by the PLUTO Collaboration [17] was based on just 111 events. However, almost all the events would have had $p_T \leqslant 3$ GeV; it is not clear to us how efficiently such events would have been triggered. On the other hand, prospects for LEP-200 certainly look very bright. Assuming that several hundred pb^{-1} will be accumulated at $\sqrt{s} = 200$ GeV (which is necessary [18] if one wants to search for Higgs bosons with masses up to M_Z), each LEP group could accumulate O(1000) events, and the p_T^γ spectrum could be studied up to $p_T^\gamma \simeq 10$ GeV.

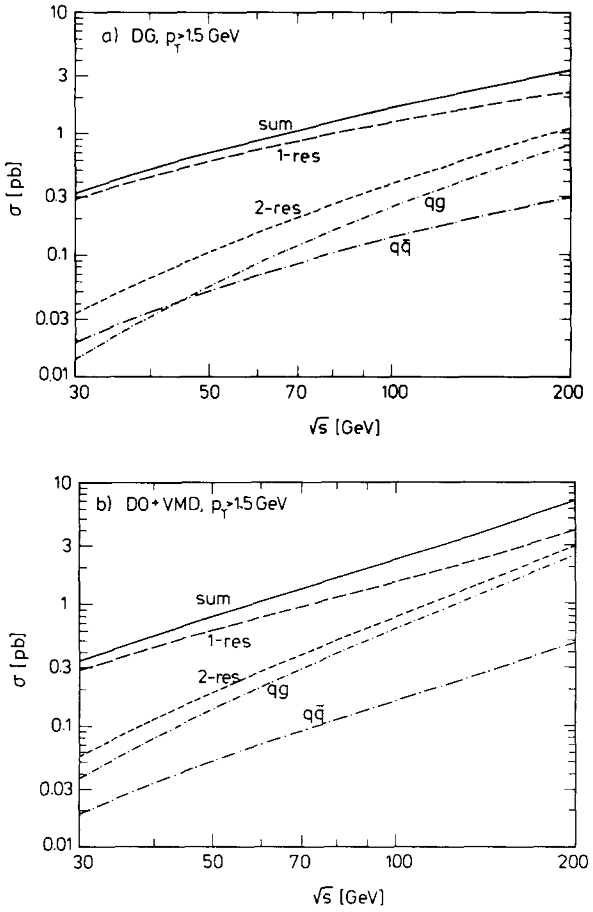

Fig. 4. The cross section for process (1) integrated over the region $p_T \leqslant 1.5$ GeV as a function of the e^+e^- centre-of-mass energy $\sqrt{s}$. Predictions using the DG and DO+VMD parametrisations are shown in (a) and (b) respectively. Parameters and notation are as in fig. 2.

Notice finally that the difference between the two parametrisations becomes more prominent at higher energies. There are two reasons for it. The double-resolved contribution increases more rapidly with $\sqrt{s}$ than the single-resolved one, due to the additional convolution with a soft distribution function, see eq. (2). In addition, for $\sqrt{s} \geqslant 80$ GeV, the single-resolved process begins to become sensitive to the region $x \leqslant 0.05$ where $q^\gamma(x)$ is not constrained by the current data. However, due to the unphysical $x \to 0$ divergence of the DO distributions, which were never intended to be used in this region, we believe that the prediction of the DG parametrisation for this contribution will be closer to the truth.

We close our discussion with some remarks concerning possible backgrounds. These can only occur if some part of the event is either not seen at all or mismeasured. Backgrounds of the first kind would be annihilation events with one or two jets going down the beam pipe, and an additional jet and a photon with roughly equal and opposite transverse momentum. Such backgrounds occur at order α^3, whereas the signal is of order α^5. However, the signal cross section also contains a factor $\sim \ln^2(s/4m_e^2) \simeq 500$ at $\sqrt{s} = 60$ GeV, and a factor $\sim \ln(p_T^2/\Lambda^2) \simeq 3$ from products of $\bar{q}^\gamma$-functions and α_s. Furthermore, to zeroth order in α_s the background can contain only two jets, one of which has to go down the beam pipe. We are interested in events where both the second jet and the photon have fairly small p_T; one can then easily convince oneself that either the jet or the photon has to emerge at large rapidity $|y| > \ln(\sqrt{s}/2p_T) = \ln(1/x_T)$. We can therefore get rid of this background, at least in principle, by simply requiring that $|y_{1,2}| < \ln(1/x_T)$, which reduces the single-resolved signal by less than 10% and effects the double-resolved signal even less. In practice one might have to tighten the cut on the jet rapidity somewhat as it might be difficult to measure it precisely at the small p_T values relevant to us. However the background will dominantly have $|y_{\rm jet}| > \ln(1/x_{\rm T})$ and $|y_\gamma| < \ln(1/x_{\rm T})$, since this is simply a back-to-back two-jet event with soft photon emission, while for the signal the two rapidities have equal distributions. In particular the configuration $y_{\rm jet} \simeq 0$, studied in fig. 3, should be almost free of this background. Note that almost half of the total signal for $p_{\rm T}^\gamma > 1.5$ GeV has $|y_{\rm jet}| \leqslant 1$.

A more dangerous class of backgrounds might arise form events with energetic jets going down each beam pipe, and a gluon "minijet" balancing the photon. Due to the softness of the minijet, the corresponding cross section would be of the same order as that for the $q\bar{q}\gamma$ events, but cannot be removed by a simple rapidity cut. Notice, however, that colour flow in this background is different from the one in the double-resolved signal events from qg-initial states, which are dominant at the TRISTAN energies and beyond. Since the presence of the (outer fringe of) the jets going down the beam pipes can presumably be detected even if the jet energy cannot be measured, this type of events would not be a background to the single-resolved signal.

An entirely different kind of background can arise from production of jets in $\gamma\gamma$ collisions. Note that the corresponding rates are [1] about 100 times higher, for given $p_{\rm T}$ and $\sqrt{s}$, than for our direct photon signal. However, at $p\bar{p}$ colliders the ratio between jets and direct photon rates is very similar; nevertheless a clean direct photon signal could be detected [19] once isolation cuts were applied. We see no reason why these cuts should be less efficient in $\gamma\gamma$ collisions.

In summary we have discussed a new class of $\gamma\gamma$ reactions, which have a hard, isolated photon in the final state. The cross section is large enough to be interesting once an integrated luminosity of O(100 $\rm pb^{-1}$) has been accumulated, which is not unusual for e^+e^- colliders. We found that at TRISTAN energies and below the single-resolved contribution dominates which is basically Compton scattering off the quarks inside the photon. This contribution is free of annihilation backgrounds once mild cuts against large rapidities are applied. We emphasize that the existing data on F_2^γ allow us to make an absolute prediction for this contribution which for $\sqrt{s} \leqslant 80$ GeV should be precise to within (20–30)%; in contrast, all the other processes involving the parton content of the photon suffer from O(100%) uncertainties mostly due to the unknown gluon content of the photon [1,2,5,7,8]. At energies higher than 80 GeV, contributions from the unknown region $x \leqslant 0.05$ become important, unless we restrict ourselves to higher $p_{\rm T}^\gamma$.

Uncertainties are again very large for our double-resolved contributions, which are dominated by the QCD Compton process $gq \to \gamma q$. The cross section for this contribution increases more rapidly with $\sqrt{s}$ than the single-resolved contribution; in addition annihilation backgrounds are more dangerous. However, these backgrounds fall with $\sqrt{s}$, so that the double-resolved contribution might be detectable at TRISTAN and beyond. Since this contribution could play a crucial role in pinning down the gluon content of the photon, in analogy with the role played by direct photon experiments at hadron accelerators [9,19], we urge our experimental colleagues to search for this contribution.

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References

[1] M. Drees and R.M. Godbole, Nucl. Phys. B 339 (1990) 355.
[2] M. Drees and R.M. Godbole, Phys. Lett. B 234 (1990) 375.
[3] For some older papers see for example, S.J. Brodsky, T.A. De Grand, J.F. Gunion and J. Weis, Phys. Rev. Lett. 41 (1978) 672; Phys. Rev. D 19 (1979) 418;
K. Kajantie and R. Raitio, Nucl. Phys. B 159 (1979) 528;
S.P. Li and H.C. Liu, Phys. Lett. B 143 (1984) 489.
[4] Ch. Berger and W. Wagner, Phys. Rep. 146 (1987) 1;
J. Olsson, Nucl. Phys. B (Proc. Suppl.) 3 (1988) 613.
[5] J.F. Owens, Phys. Rev. D 21 (1979) 54;
P. Aurenche, R. Baier, A. Douiri, M. Fontannaz and D. Schiff, Nucl. Phys. B 286 (1987) 553;
H. Baer, J. Ohnemus and J.F. Owens, Phys. Rev. D 40 (1989) 2844.
[6] D.W. Duke and J.F. Owens, Phys. Rev. D 26 (1982) 1600.
[7] M. Drees and R.M. Godbole, Phys. Rev. Lett. 61 (1988) 682; Phys. Rev. D 39 (1989) 169;
W.J. Stirling and Z. Kunszt, in: Proc. HERA Workshop, ed. R.D. Peccei, p. 331;
H. Baer, J. Ohnemus and J. Owens, Z. Phys. C 42 (1990) 657;
A.C. Bawa, M. Krawczyk and W.J. Stirling, DESY preprint 90-123 (1990).

[8] M. Drees and F. Halzen, Phys. Rev. Lett. 61 (1988) 275;
R. Gandhi, I. Sarcevic, A. Burrows, L. Durand and H. Pi, Phys. Rev. D 43 (1990) 263;
T.K. Gaisser, F. Halzen, T. Stanev and E. Zas, Phys. Lett. B 243 (1990) 444.
[9] J.F. Owens, Rev. Mod. Phys. 59 (1987) 465.
[10] E. Witten, Nucl. Phys. B 120 (1977) 189.
[11] C.F. von Weizsäcker, Z. Phys. 88 (1934) 612;
E.J. Williams, Phys. Rev. 45 (1934) 729.
[12] S.J. Brodsky, T. Kinoshita and H. Terazawa, Phys. Rev. D 24 (1971) 1532.
[13] R. Bhattacharya, J. Smith and G. Grammer, Phys. Rev. D 15 (1977) 3267.
[14] M. Drees and K. Grassie, Z. Phys. C 28 (1985) 451.
[15] PLUTO Collab., Ch. Berger et al., Phys. Lett. B 142 (1984) 111.
[16] M. Glück, K. Grassie and E. Reya, Phys. Rev. D 30 (1984) 1447.
[17] PLUTO Collab., Ch. Berger et al., Phys. Lett. B 107 (1981) 164.
[18] Z. Kunszt and W.J. Stirling, Phys. Lett. B 242 (1990) 507;
N. Brown, Rutherford Laboratory preprint RAL-90-059 (1990).
[19] For a recent review see e.g. J. Huston, Proc. 1989 Intern. Symp. on Lepton and photon interactions at high energies, ed. M. Riordan (World Scientific, Singapore, 1990) p. 348.