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Minijets and large hadronic backgrounds at e^+e^- supercolliders.

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Abstract

We show that the hadronic structure of the photon, along with the beamstrahlung phenomenon, leads to a very large rate of $\gamma\gamma \to$ jets at e^+e^- supercolliders. Already at $\sqrt{s} = 1$ TeV, for round, dense bunches, we expect O(5) to O(50) 'minijet' events per bunch crossing, giving rise to an 'underlying event'. Thus e^+e^- supercolliders will be 'messier' than expected, unless beamstrahlung can be kept under control.

In recent years much effort has been devoted [1, 2] to the exploration of the physics capabilities and design problems of future e^+e^- (super) colliders. The main argument in favour of e^+e^- machines as compared to pp supercolliders [1, 3] like the SSC or LHC is the supposed 'cleanliness' of the electron machines, i.e. the absence of an 'underlying event' produced by the spectator jets which are part of any hard interaction event at hadron colliders, spraying a multitude of usually soft hadrons over the detector, which cause many background problems. The purpose of this letter is to point out that, due to the hadronic structure which (quasi-)real photons develop to $O(\alpha)$ [4], photon-photon interaction rates at future e^+e^- colliders will be very large, making such machines much 'messier' than usually anticipated. This problem is aggravated by the beamstrahlung phenomenon [5] , which already at center-of-mass energies $\sqrt{s}$ around 1 TeV can increase the effective single photon flux by an order of magnitude. These large two-photon rates cause two distinct problems:
i) A large majority of all 'hard' events (defined as depositing more than some minimal amount of transverse energy, $E_{T,min}$, in the central part of the detector) could be two-photon events, potentially obscuring signals for 'new physics'. This can already happen at LEP 200 energies and beyond, but can be alleviated by increasing $E_{T,min}$ with only a moderate loss of signal, unless beamstrahlung is very large.
ii) For supercollider energies ($\sqrt{s} \simeq 1$ TeV) the **total** rate for $e^+e^- \to e^+e^-$ + jets, including 'minijets' [6] with transverse momentum $p_T \simeq 1 - 3$ GeV, can become so large that one or several of these reactions occur at **every** bunch crossing. These reactions will then effectively produce an 'underlying event'; for $\sqrt{s} \geq 2$ TeV, this underlying event could contain as much **transverse** energy as a single annihilation event, making the situation **worse** than at pp supercolliders, unless beamstrahlung can be kept under control.

It has been noted [7] that beamstrahlung can cause yet another problem by giving rise to large rates for e^+e^- pair production; already at $\sqrt{s} = 500$ GeV , there could be as many as $O(10^6)$ pairs per bunch crossing [8]. However here the typical transverse momenta will be of $O(m_e)$, so that most of the electrons will be either boosted out of the detector or curl up in the magnetic field. Requiring that these electrons have $p_T > 200$ MeV with rapidities less than 2, we estimate that only 1 of these hard events will survive, if we assume that coherent production no longer contributes at these large transverse momenta. This gives a negligible contribution to the total transverse energy, which moreover will be fully contained in the electromagnetic calorimeter. In contrast the spectator hadrons naturally have $p_T \simeq 300$ MeV, and are thus expected to penetrate in the hadronic calorimeters.

For the purpose of this letter we need only discuss the inclusive production of two high-p_T jets in the collision of two (quasi-)real photons. Three different classes of processes [9] contribute to this reaction (see fig. 1). The 'direct process' of fig. 1a is due to the simple $\gamma\gamma \to q\bar{q}$ production, present already in the naive quark-parton model. However, to first order in α the photon also develops [4] a non-vanishing quark and gluon content. It is thus possible to 'pull' quarks or gluons out of (quasi -)real photons in much the same way they can be 'pulled out' of nucleons. Fig. 1b depicts the case where only one photon is resolved into its partonic components, which then

interact with the other photon; we call these the 'single-resolved' processes ('1-res' for short). Finally, fig. 1c shows the situation where both photons are resolved, so that the hard scattering is a pure QCD $2 \to 2$ process; we call these the 'twice-resolved' contributions ('2-res' for short). It is very important to note here that every resolved photon will produce [10] a spectator jet of hadrons with small transverse momentum relative to the initial photon direction, which for (quasi-)real photons coincides with the electron (or positron) beam direction.

The cross section for the twice-resolved contributions can be calculated from

$$\frac{d\sigma}{dp_T\, dy_1 dy_2} = 2p_T\, x_m\, z_m \int_{x_m}^1 \frac{dx}{x} f_{\gamma/e}(\bar{E},x)\, \bar{q}^\gamma(\frac{x_m}{x},Q^2) \int_{z_m}^1 \frac{dz}{z} f_{\gamma/e}(\bar{E},z)\, \bar{q}^\gamma(\frac{z_m}{z},Q^2)\, \frac{d\hat{\sigma}}{d\hat{t}}(\hat{s},\hat{t},\hat{u}), \quad (1)$$

where

$$\begin{aligned} x_m &= \frac{x_T}{2}(e^{y_1}+e^{y_2}); \\ z_m &= \frac{x_T}{2}(e^{-y_1}+e^{-y_2}); \\ x_T &= 2p_T\ /\sqrt{s}\ ; \\ \hat{s} &= s\, x_m\, z_m\ . \end{aligned}$$

The kinematic boundaries for the two jet rapidites y_1 and y_2 are

$$|y_1| < \mathrm{arcosh}(\frac{1}{x_T}); \qquad -\ln(\frac{2}{x_T} - \exp(-y_1)) < y_2 < \ln(\frac{2}{x_T} - \exp(y_1)). \quad (2)$$

The hard sub cross sections [11] $\hat{\sigma}$ have to be taken at the point

$$\hat{t} = -\frac{\hat{s}}{2}(1+\sqrt{1-4p_T^2/\hat{s}}), \quad (3)$$

and to be symmetrized under $\hat{t} \leftrightarrow \hat{u}$ since we do not distinguish between jets of different origin (q,$\bar{q}$ or G). For the single-resolved (direct) processes, one (both) of the parton density functions $\bar{q}^\gamma(x,Q^2)$ have to be replaced by $\delta(1-x)$, and the proper hard sub-process cross sections have to be inserted [12]. Notice that all the three classes of contribution are of the same order in coupling constants, since $\bar{q}^\gamma(x,Q^2)$ is $O(\alpha/\alpha_s)$ [4].

Our choices for the parton densities inside the photon $\bar{q}^\gamma(x,Q^2)$ and the photon fluxes inside the electron $f_{\gamma/e}(\bar{E},z)$ require some discussion. At present only

two parametrizations of the $\bar{q}^\gamma(x,Q^2)$ exist. The parametrization of [12] uses the 'asymptotic' prediction of Witten [4]. For our choice $\Lambda_{QCD} = 400$ MeV, this has to be augmented by a 'hadronic' component, which can be estimated from the vector meson dominance (VMD) model, in order to describe data [13] on the electromagnetic structure function F_2^γ of the photon (see ref.10 for more details). However, this parametrization cannot (and was never intended to) be used at very small x, since here it diverges even worse than the $x^{-1.6}$ behaviour of the exact 'asymptotic' prediction for $\bar{q}^\gamma(x,Q^2)$. Since the small-x region will obviously play an increasingly important role at higher energies, we have modified the original parametrization of [12] in the region $x < 0.05$:

$$q_i^\gamma(x,Q^2) = c_i \ln(Q^2/\Lambda^2)\ x^{-1.6} \quad (4)$$

where the c_i are fixed to give a smooth transition at $x = 0.05$; the same $x^{-1.6}$ behaviour has been assumed for the gluon density at small x. Note that the VMD contribution [10] is not affected by this modification, as it is well-behaved for $x \to 0$. We henceforth call this the 'modified DO + VMD' parametrization.

While the $x^{-0.6}$ behaviour of the 'asymptotic' leading-order prediction for F_2^γ may still be acceptable since the counting rates remain finite, it signals a fundamental problem of the 'asymptotic' prediction, which leads [14] to an x^{-1} divergence with a negative sign in second order, and even worse $x \to 0$ poles in yet higher orders [15]. The 'DG' parametrization [16] avoids this problem, since it parametrizes the well-behaved [17] Q^2 evolution of some input distributions at scale $Q_0^2 = 1$ (GeV)^2, which were chosen such that a preliminary version of the data of [13] were reproduced; as shown in ref. 10 it also fits the final version of these data quite well.

From a purely phenomenological point of view the use of the DG parametrization probably leads to a conservative estimate of two-photon cross sections, since here it is assumed that gluons are only created radiatively inside the photon. In contrast, due to their very rapid increase at low x (see eq.(4)) and the large VMD contribution in particular to the gluon content of the photon, the use of the modified DO+VMD parametrization might well overestimate two-photon cross sections. In all cases we have chosen the momentum scale $Q^2 = \hat{s}/4$, both in $\bar{q}^\gamma(x,Q^2)$ and in the running QCD coupling α_s; furthermore, we have varied the number N_f of active flavors with Q^2, using $N_f = 3$ for $Q^2 < 50$ (GeV)^2 and $N_f = 5$ for $Q^2 > 500$ (GeV)^2. This should be a conservative treatment since we have, e.g., set the charm contribution to zero for $\hat{s} < 200\ (GeV)^2$.

Finally, our choice for the photon flux function $f_{\gamma/e}$ needs to be discussed. For the flux of photons which are not resolved into their partonic constituents we have used the expression which includes contributions beyond the leading logarithmic terms [18]. Notice that this expression has been derived by integrating the virtuality P^2 of the exchanged photon over the full kinematic range; in particular, the contribution of far off-shell photons is suppressed only by the $1/P^2$ photon propagator. On the other hand, it is known [19] that the parton content of the photon decreases once $P^2 >> \Lambda^2$. Indeed, if the virtuality P^2 of the photon is larger than the scale Q^2 at

which it is probed, the concept of photon structure function becomes meaningless. In this region the whole process can be described fully perturbatively, without having to resort to structure functions. To be conservative, we have ignored these contributions altogether. Furthermore, in order to take into account the suppression of virtual photon structure function for $\Lambda^2 < P^2 < Q^2$, we have multiplied the flux of resolved photons with a suppression factor of 0.85, which has been estimated from numerical results of Rossi [19]. Altogether we thus have for the effective flux of resolved photons due to bremsstrahlung:

$$f_{\gamma/e}^{res}(z) = 0.85\frac{\alpha}{2\pi}\frac{(1+(1-z)^2)}{z}\ln(\frac{Q^2}{m_e^2}). \qquad (5)$$

where m_e is the electron mass and Q^2 the scale at which the photon is probed.

We are now in a position to present results on two-photon production of jets in the case when beamstrahlung is negligible. This is expected to be true for $\sqrt{s} < 500$ GeV. (However already at the planned JLC ($\sqrt{s} = 500$ GeV) beamstrahlung effects could be sizable [20].) In fig. 2 we show the cross section (1) integrated over the region $|y_{1,2}| < 2$, $p_T > 5$ GeV, as a function of $\sqrt{s}$. The short dashed curve shows the contribution of the direct $\gamma\gamma \to q\bar{q}$ production mechanism. For the chosen kinematical cuts, which are meant to guarantee the occurence of two well-defined jets in the central region of the detector, it still dominates at LEP 200 energies, $\sqrt{s} \leq 200$ GeV. However, it increases only logarithmically with the energy, while both the single resolved (long-short dashed) and twice-resolved (long dashed) contributions increase much faster [10] with $\sqrt{s}$, almost linearly in the latter case. The solid curve shows the prediction for the total two jet rate using the DG parametrization. Notice that even for our cuts, which are quite stringent for two photon physics, the direct contribution alone would underestimate the total rate by at least a factor of 5 already at the intermediate energy $\sqrt{s} = 500$ GeV, which is now envisioned for the next generation of e^+e^- machines [2]. If the photon structure is better described by the modified DO+VMD parametrization (dot-dashed curve), this factor is closer to ten.

Notice also that the cross sections of fig. 2 are huge compared to typical annihilation cross sections. For instance, the QED cross section for $e^+e^- \to \mu^+\mu^-$ amounts to only 0.4 pb at $\sqrt{s} = 500$ GeV, less than 0.3% of our hard two-photon cross section; most cross sections for the production of new, heavy particles will be even smaller. Nevertheless, at these energies the two-photon background can still be brought under control quite easily; it can be reduced to the level of the annihilation cross sections by increasing the p_T-cut to about 20 GeV, which would only have mild effects on most 'new physics' searches. (An exception is the search for final states containing heavy stable neutral particles.) On the other hand, the difference between the solid and dot-dashed curve indicates that a 500 GeV machine would be very well suited to study $\gamma\gamma$ reactions in detail. (Of course, prior to that our knowledge of the hadronic structure of the photon will have increased due to two-photon data from TRISTAN [10] and LEP [21], as well as photoproduction data from HERA [22].)

The situation will be quite different, however, for $\sqrt{s} \geq 1$ TeV. We see from fig. 2 that the ratio of two-photon events from bremsstrahlung photons to annihilation

events increases roughly like $s^{3/2}$. In addition, at these energies beamstrahlung [5, 23] can become very important. It is due to the coherent interaction of the particles in one bunch with a large number of particles in the other bunch, and therefore critically depends on the density of the bunches, which in turn is related to the luminosity per bunch crossing. Notice that beamstrahlung produces truly real (on-shell) photons, so that the hadronic structure of the photon is expected to play an even greater role than for photons produced by bremsstrahlung.

In order to estimate how big the effect of beamstrahlung could be at $\sqrt{s} = 1$ TeV we have parametrized the beamstrahlung contribution to $f_{\gamma/e}$, using the most extreme result of [23], which assumes round beams, a bunch length of 0.3 mm, and a luminosity of $2.8\text{x}10^{30}\text{cm}^{-2}$ per bunch crossing. Their result can be written as

$$f_{\gamma/e}^{beam} = \begin{cases} (2.25 - \sqrt{\frac{x}{.166}})(\frac{1-x}{x})^{2/3} & x < 0.84 \\ \\ 0 & x > 0.84. \end{cases} \qquad (6)$$

Of course, the bremsstrahlung contribution has to be added to this [24]

In fig. 3 we show the resulting p_T spectrum as predicted by the DG parametrization. The upper four curves use the same notation as in fig. 2 and include the beamstrahlung contribution according to eq.(6); the lowest (dot-dashed) curve shows the total result for the case without beamstrahlung. We see that if beamstrahlung is as strong as assumed in eq.(6), it increases the two-photon event rate by at least two orders of magnitude over the whole range of p_T values shown. As a result, one would have to increase the p_T-cut to around 120 GeV in order to reduce the two-photon cross section to the level of typical annihilation cross sections. In view of the fact that in the presence of strong beamstrahlung many annihilation events contain considerably less energy than the nominal $\sqrt{s}$ of the machine, this is a significant cut, in particular if the signal cross section is also peaked at small angles, which is true, e.g., for the $e^+e^- \to VV(V = W, Z)$ cross sections.

So far we have restricted our discussion to events with relatively hard, central jets. We see that even with these cuts we get roughly one event with $p_T > 5$ GeV every 10 bunch crossings, if beamstrahlung is as strong as assumed in eq.(6). This naturally brings to our mind the question how large these event rates become if we relax our cuts to the limits of applicability of perturbative QCD. The result is shown in fig. 4, where we give the cross section (1) integrated over the full kinematically allowed range (2) of rapidities and have varied the lower limit of the p_T integration, for fixed $\sqrt{s} = 1$ TeV. The solid (dashed) curves are with maximal (no) beamstrahlung; for a given pattern, the upper (lower) curve is for the modified DO+VMD (DG) parametrization. All curves sum direct, single-resolved and twice-resolved contributions, but we see from figs. 2 and 3 that the twice-resolved contributiuons are by far the dominant ones.

Obviously our result depends strongly on $p_{T,min}$. Unfortunately, this parameter cannot be predicted from first principles. Analyses of pp and γp scattering might, however, serve as indications. The idea that QCD jets with p_T in the GeV range could drive the observed rise in total hadronic cross sections dates back more than

15 years [6]. More recent analyses [25] show that the rise of the total pp and $\bar{p}p$ cross sections can be well described by minijets with $p_{T,min} \simeq 1.3$ - 2 GeV, for $\sqrt{s} \leq 1$ TeV. Similar values of $p_{T,min}$ have recently been claimed [22, 26] to describe the energy dependence of the total γp cross section, although here a value of $p_{T,min}$ as low as 1 GeV can be brought into agreement with the data [27]. In view of these results the choice of the DG parametrization with $p_{T,min} \simeq 2 - 2.5$ GeV seems quite conservative. We see from fig. 4 that for the case of maximal beamstrahlung, this translates into 4-8 minijet events per bunch crossing! Allowing for somewhat smaller values of $p_{T,min}$ and/or a larger gluon content of the photon, we can at present not exclude the possibility that the true number is an order of magnitude higher than this.

What are the consequences of these minijet reactions, which will accompagny all annihilation events? Recall that the twice-resolved contribution, which is by far the dominant one, is characterized by two spectator jets. For the case of fig. 4 we find that these spectator jets have typical energies around 100 GeV. Of course, most of this energy will go in the very forward direction and thus might not be detected at all. Notice, however, that there is nontrivial color flow between the spectator jets and the hard (mini-)jets. One can therefore expect the spectator jets to deposit 1.5-2 GeV transverse energy per unit of rapidity. Adding 4 GeV transverse energy from the hard jets and integrating over the region $|y| < 2$ we thus estimate that each minijet event will deposit about 10 GeV transverse energy in the central part of the detector. For 10 simultaneous events this yields an 'underlying event' contributing as much as 100 GeV transverse energy which is considerably worse than at hadron supercolliders. We emphasize again that, except for the choice of maximal beamstrahlung (6), we consider our estimates, of both the event rate and the energy deposited by each event, as quite conservative.

Obviously, the traditional 'cleanliness' of e^+e^- colliders can only be preserved if beamstrahlung can be reduced considerably. As discussed in ref.23, this can be achieved by using elliptical beams with large ratio r of semimajor and semiminor axes. However, if the values of the bunch length and luminosity per bunch crossing are left unaltered, one would have to chose r as large as 100 in order to reduce the beamstrahlung contribution to $f_{\gamma/e}$ roughly to the level of the bremsstrahlung contribution, leading to a minijet cross section roughly 4 times bigger than indicated by the dashed curves in fig. 4. This could still take us dangerously close to the level of one minijet event per bunch crossing. Moreover, making the beams very flat will obviously make it even more difficult to achieve head-on collisions. Furthermore, the number of minijet events increases very rapidly with $\sqrt{s}$. Obviously, the machine luminosity has to be increased like s in order to achieve a constant rate of annihilation events. This will lead to an even more rapid increase of the minijet rate, since a larger luminosity will necessitate denser bunches and hence more beamstrahlung per electron. Finally, the average photon energies will also increase with $\sqrt{s}$; recall that minijet cross sections increase with the center-of-mass energy of the parent particles. It thus seems very likely to us that for $\sqrt{s} \geq 2$ TeV the existence of an 'underlying event' with large transverse energy content can only be avoided if the number of bunches can somehow be increased considerably, which would allow for less

dense bunches without reducing the integrated annihilation event rate. Most recent designs [28] for machines with $\sqrt{s} = 1$ TeV indeed use flat beams and a number of bunches that is larger than the 'cycle rate' of the accelerator. At those machines the hadronic background due to beamstrahlung should therefore be smaller than in our worst case scenario; however, it can still be very serious. We think that this problem should be considered when plans for the next generation of accelerators are finalized.

Clearly the validity of our conclusions crucially hinges on the existence of the resolved processes of figs. 1b, 1c. Using the traditional VDM picture [29] with a constant total $\gamma\gamma \to$ hadrons cross section of about 250 nb also leads to large two-photon cross sections of 1.9 (0.4) μb at $\sqrt{s} = 1$ TeV, for $\gamma\gamma$ center-of-mass energy $W > 10$ (100) GeV, if the beamstrahlung specturm of eq.(6) is assumed. However, these are diffractive events which deposit very little energy in the central part of the detector. In this context it is interesting to note that models based on the incoherent sum of these soft contributions and the direct process of fig. 1a have consistently failed to describe data [30] on jet production in the collision of (quasi-)real photons at PEP and PETRA. The missing 'third component' shows a softer p_T spectrum and broader thrust distribution than the direct contribution, as expected [10] from the resolved contributions. Since the relative importance of these contributions increases with energy, data from the TRISTAN collider should help to test the overall validity of our picture. Note also that, if the resolved contributions are to be identified with the 'third component', they must manifest themselves already at $p_T \simeq 2 - 3$ GeV, since in this region the excess of the data [30] over the prediction of the 'direct+soft' models is largest.

In summary we have pointed out that the combination of the beamstrahlung phenomenon with QCD predictions for the hadronic structure of the photon leads to very large two-photon event rates at e^+e^- supercolliders. Already at $\sqrt{s} = 1$TeV the two-photon contribution can dominate total two-jet production up to $p_T \simeq 100 - 130$ GeV if round, dense bunches are used in the accelerator. With such bunches one also has to expect O(5), or possibly even O(50), minijet events per bunch crossing, leading to an 'underlying event'. Data from TRISTAN and HERA should help to sharpen these predictions, which are obviously relevant for the planning of future e^+e^- accelerators and experiments.

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Figure captions :

fig. 1 Typical Feynman diagrams for the direct (1a), once-resolved (1b) and twice-resolved (1c) contributions to the production of high-p_T jets in $\gamma\gamma$ collisions at e^+e^- colliders. Note that each resolved photon produces a spectator jet in addition to the hard jets.

fig. 2 The cross section of eq.(1) integrated over the region $|y_{1,2}| \leq 2$, $p_T > 5$GeV. The dot-dashed curve shows the total cross section when the DO parametrization [12], modified as described in the text and augmented by a VMD inspired part [10], is used, while the short dashed curve shows the contribution from the direct contribution alone. The long-short dashed, long dashed and solid curves show the once-resolved, twice-resolved and total contributions, respectively, for the DG parametrisation [16]. We have used $Q^2 = \hat{s}/4$ and a floating number of active flavors, as described in the text. The effects of beamstrahlung are not included.

fig. 3 The cross section (1) integrated over $|y_{1,2}| \leq 2$, $\sqrt{s} = 1$ TeV as a function of the jet transverse momentum p_T. The upper four curves include the beamstrahlung contribution according to eq.(6) and use the same notation as fig.2, while the lowest (dot-dashed) curve shows the total result if beamstrahlung is neglected. All results are for DG parametrization [16]. The kinks that occur in the curves at $p_T \simeq 7$ GeV are due to the fact that we have conservatively switched off the charm contribution for $Q^2 < 50$ (GeV)^2, as described in the text.

fig. 4 The cross section (1) integrated over the whole kinematically allowed range of rapidities and over $p_T > p_{T,min}$, at $\sqrt{s} = 1$ TeV, with (solid) and without (dashed) the beamstrahlung contribution of eq.(6). The upper (lower) curves of given pattern are for the modified DO+VMD (DG) parametrization. All curves are for the total $\gamma\gamma$ cross section; however, the single-resolved and direct contributions amount to only O(10%) of the total.

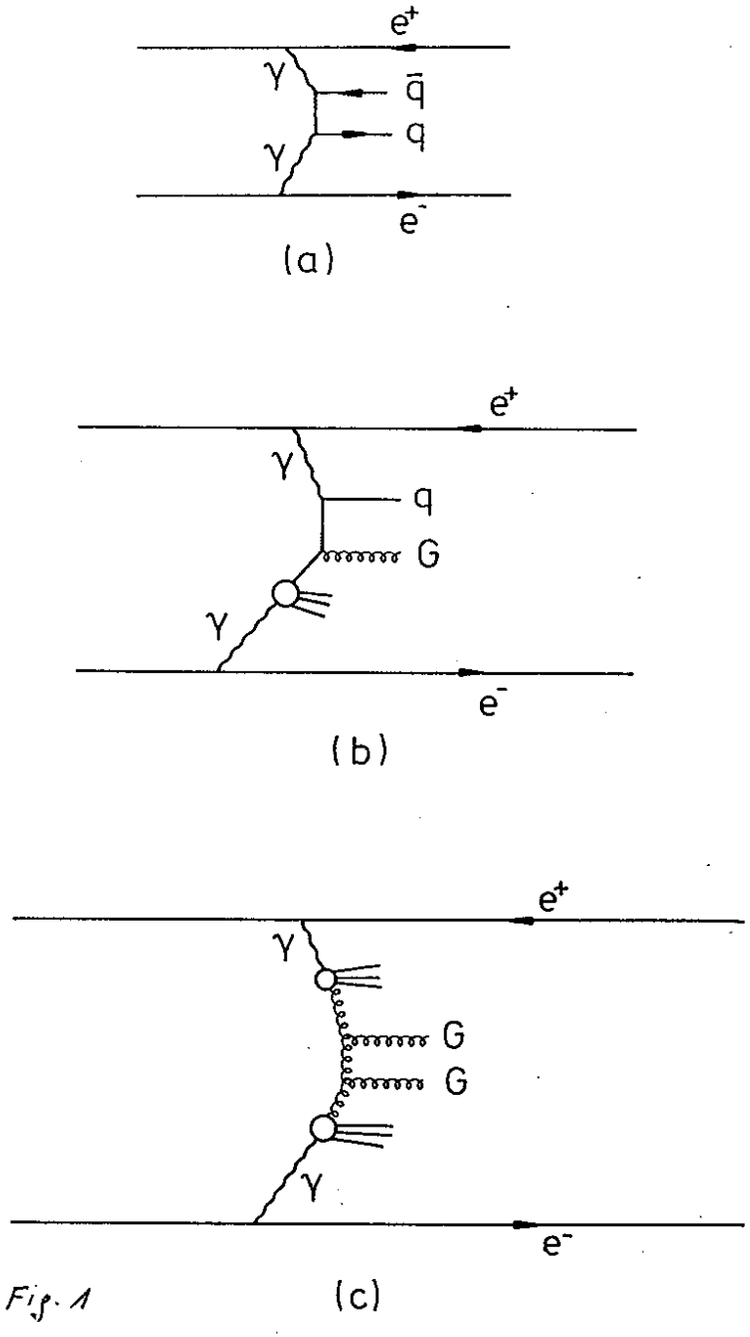

Fig. 1

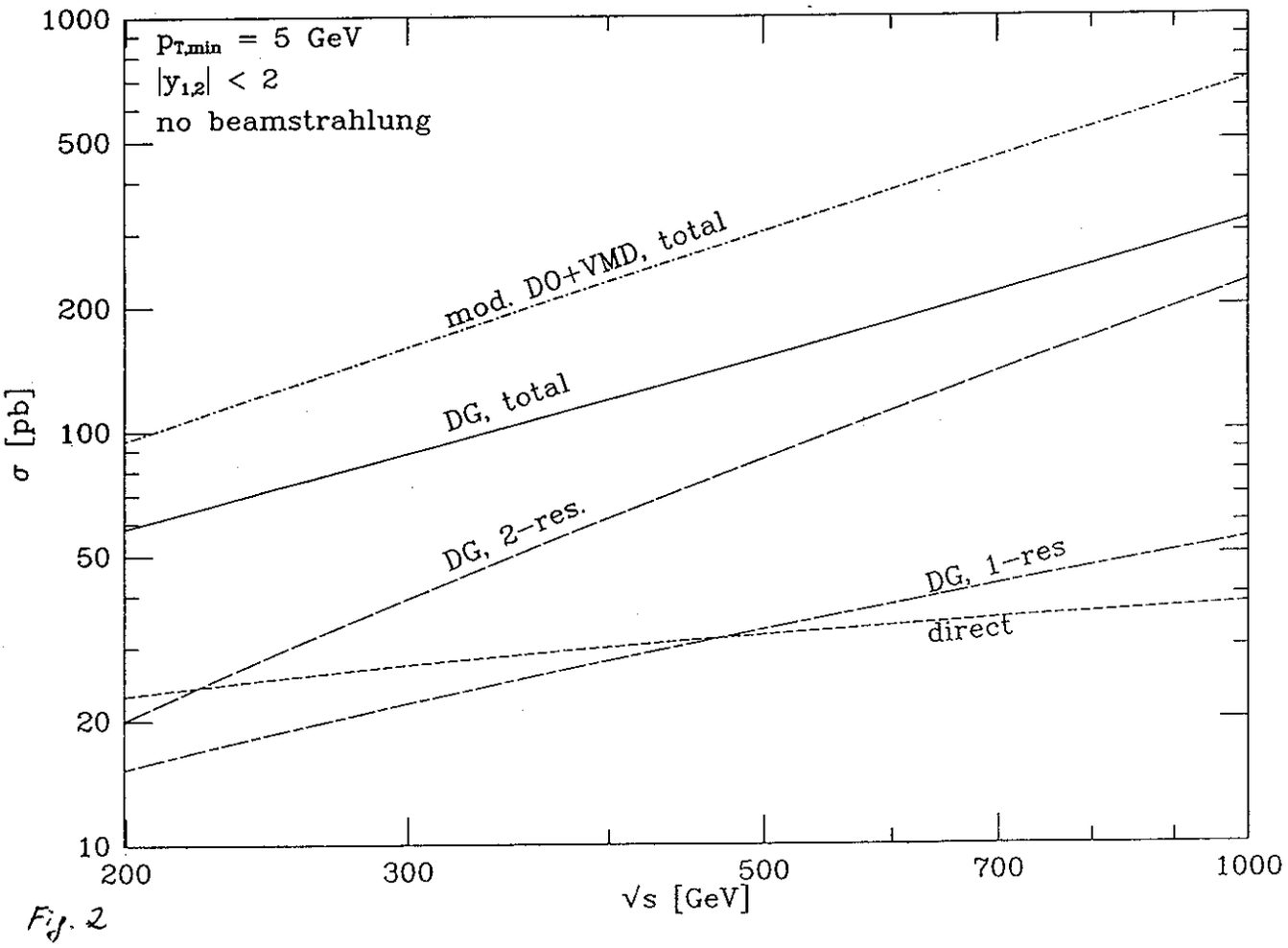
$p_{T,min}$ = 5 GeV
$|y_{1,2}| < 2$
no beamstrahlung
mod. DO+VMD, total
DG, total
DG, 2-res.
DG, 1-res
direct
σ [pb]
√s [GeV]
1000
500
200
100
50
20
10
200
300
500
700
1000

Fig. 2

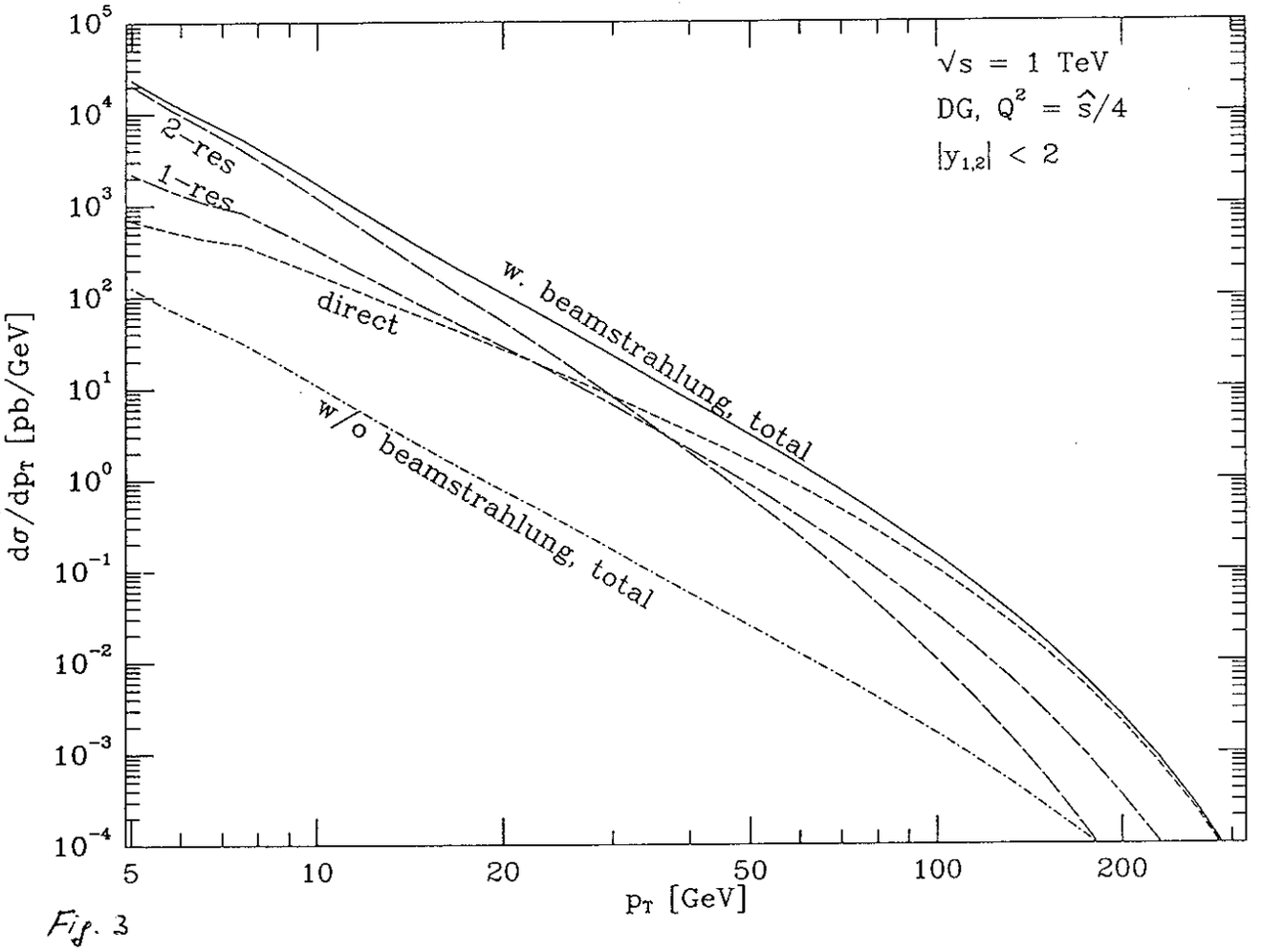
√s = 1 TeV
DG, $Q^2 = \hat{s}/4$
$|y_{1,2}| < 2$
2-res
1-res
direct
w. beamstrahlung, total
w/o beamstrahlung, total
$d\sigma/dp_T$ [pb/GeV]
p_T [GeV]
10^5
10^4
10^3
10^2
10^1
10^0
10^{-1}
10^{-2}
10^{-3}
10^{-4}
5
10
20
50
100
200

Fig. 3

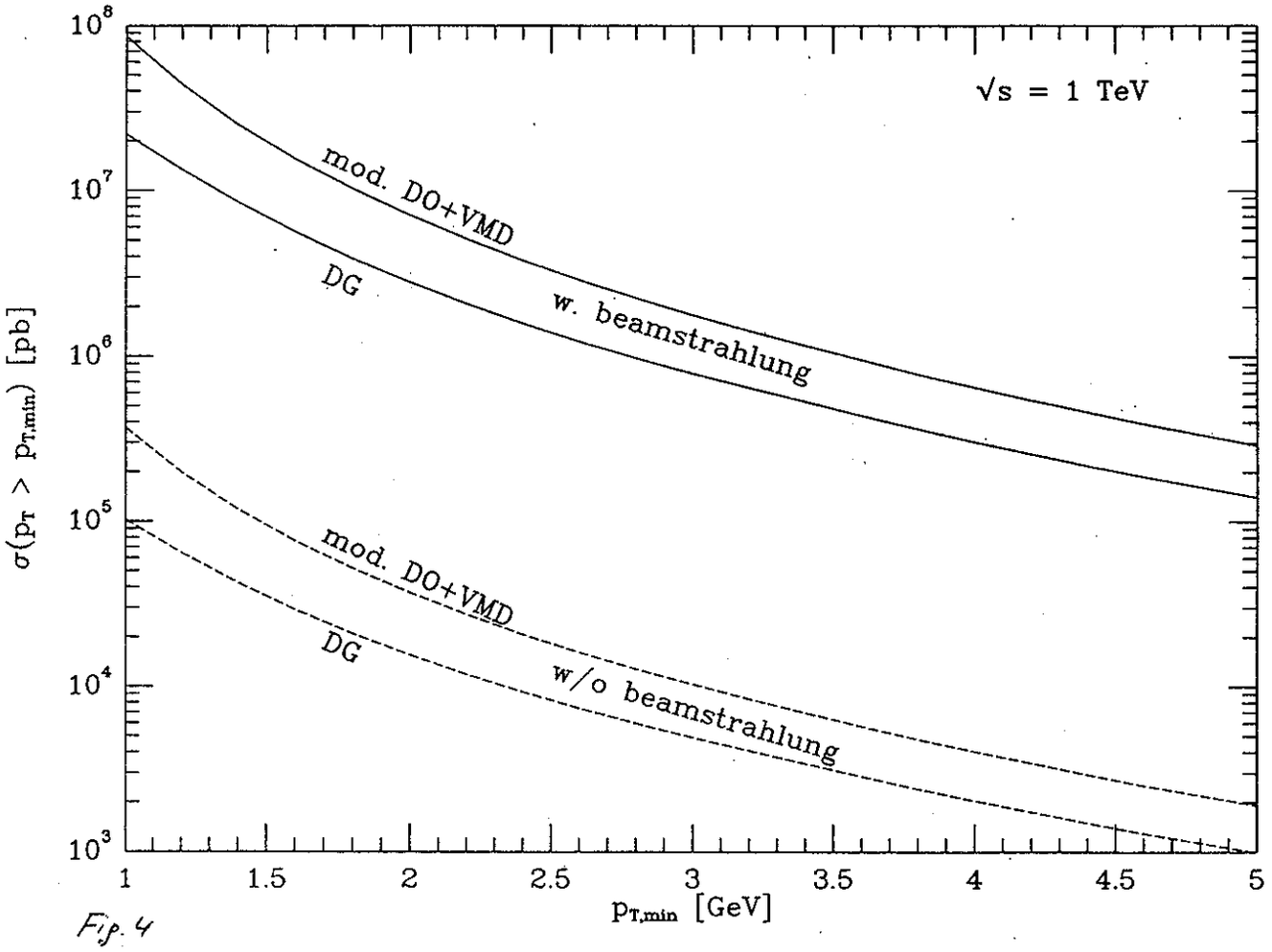
√s = 1 TeV
mod. DO+VMD
w. beamstrahlung
DG
mod. DO+VMD
w/o beamstrahlung
DG
σ(p_T > p_T,min) [pb]
p_T,min [GeV]
10^8
10^7
10^6
10^5
10^4
10^3
1
1.5
2
2.5
3
3.5
4
4.5
5

Fig. 4