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BEAMSTRAHLUNG AND THE QED, QCD BACKGROUNDS IN LINEAR COLIIDERS*

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ABSTRACT

The intense radiation, called *beamstrahlung*, during the collision of e^+e^- beams in a linear collider, is reviewed, with attention to the influence of beam-beam disruption on the beamstrahlung spectrum. We then discuss the various detector backgrounds induced by these hard beamstrahlung photons, as well as the Weizsacker-Williams photons, through various QED and QCD processes, namely the coherent and incoherent e^+e^- pair creation and the hadron production and minijet yields.

1. Introduction

One of the most important issues in the design of future e^+e^- colliders is the effect of the beam-beam interaction on the physics environment. The single-pass nature of linear colliders necessitates the need for colliding tiny, intense bunches of electrons and positrons in order to achieve the required high luminosity. In this circumstance, these bunches interact strongly with one another, producing large numbers of hard photons, a phenomenon called *beamstrahlung*.[1] This effect potentially creates troublesome backgrounds for experiments on e^+e^- annihilation and must be controlled by adjustment of the collider parameters or the interaction region geometry.

Earlier, Zolotarev et al.[2] studied the e^+e^- pair creation backgrounds from the collision of beamstrahlung photon and the individual particle in the oncoming beam. Chen and Telnov[3] first pointed out that there is a very high probability for the beamstrahlung photons to turn into e^+e^- pairs through *the coherent* interaction between the photon and the collection of the opposing bunch particles. Beyond a certain threshold, a large fraction of beamstrahlung photons will turn into such pairs.[3,4] Recently, Drees and Godbole[5,6] called attention to another potentially serious background due to the beam-beam interaction: They proposed that photons created by the bunch collision can interact to produce hadronic jets. In some designs, the rate of this process exceeds one jet pair per bunch crossing. Under these conditions, each e^+e^- annihilation event would be superposed on an extraneous system of hadronic jets. Further investigations into this issue, however, suggest a somewhat lower estimate on the minijet cross section.[7,8]

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In this paper, we will first review the beamstrahlung spectrum, with attention to the effective beamstrahlung due the beam deformation during beam-beam collision. We then turn to the coherent and the incoherent e^+e^- pair creation processes from both beamstrahlung and bremsstrahlung photons, in Section 3. We will show that while the coherent pair production may be more abundant beyond certain threshold, there is nevertheless a way to stay below this threshold by properly adjusting the beam parameters. On the other hand, the incoherent pairs with inherently large angles can not be avoided. In Section 4, we discuss the hadron production and the minjet problem. We review the key ingredients in the so-called Reference Model introduced in Ref. 8., and compare it with the Drees-Godbole minijet model. The various backgrounds are then estimated for the next generation linear colliders currently under study.

2. Beamstrahlung Spectrum

In contrast to bremsstrahlung, beamstrahlung occurs in the situation where the scattering amplitudes between the radiating particle and the target particles within the characteristic length add coherently. Typically for the beam-beam collision in linear colliders there can be well over a million target particles involved within the coherence length. The process can therefore be well described in a semi-classical calculation where the target particles are replaced by their collective EM fields.

High energy e^+e^- beams generally follow Gaussian distributions in the three spatial dimensions, and their local field strength varies inside the beam volume. In the weak disruption limit, where particle motions are para-axial, it is possible to integrate the radiation process over this volume and derive relation which depend only on averaged, global beam parameters.[9] The overall beamstrahlung intensity is controlled by a global *beamstrahlung parameter*,

$$\Upsilon_0 = \gamma\frac{\langle B\rangle}{B_c} = \frac{5}{6}\frac{r_e^2\gamma N}{\alpha\sigma_z(\sigma_x+\sigma_y)} \quad , \tag{2.1}$$

where $\langle B\rangle$ is the mean electromagnatic field strength of the beam, $B_c = m_e^2/e \simeq 4.4\times10^{13}$ Gauss is the Schwinger critical field, N is the total number of particles in a bunch, $\sigma_x, \sigma_y, \sigma_z$ are the nominal sizes of the Gaussian beam, γ is the Lorentz factor of the radiating particle, r_e is the classical electron radius, and α is the fine structure constant.

The collective fields in the beam also deform the other beam during collision, by an amount controlled by a global *disruption parameter*:[10]

$$D_{x,y} = \frac{2Nr_e\sigma_z}{\gamma\sigma_{x,y}(\sigma_x+\sigma_y)} \quad . \tag{2.2}$$

In the most general designs for linear colliders, the photon spectrum due to beamstrahlung is not a factorized function of the electron and positron sources and depends on the detailed evolution of the bunches in the collision process. In general, then, the spectrum of radiation must be computed by detailed simulation. However, typical beams in linear colliders are very long and narrow. Since all particles oscillate within the focusing potential that is defined by the geometry of the oncoming beam, the oscillation amplitudes are small compared with its periodicity. To this end the para-axial assumption of particle motion is still approximately valid. Then the main effect of disruption on beamstrahlung is the change of effective EM fields in the bunch due to the deformation of the transverse beam sizes. Thus beamstrahlung is in practice still factorizable even under a non-negligible disruption effect, if only an effective beam size can be derived.

To find the effective beam size, we resort to the so-called luminosity enhancement factor, defined as the ratio of the effective luminosity to the nominal luminosity, due to the change of beam size:

$$H_D \equiv \frac{L}{L_0} = \frac{\sigma_x \sigma_y}{\bar{\sigma}_x \bar{\sigma}_y} \quad , \tag{2.3}$$

The *luminosity enhancement factor is calculable analytically only in the* $D \ll 1$ limit. Beyond this limit the dynamics of beam-beam interaction becomes nonlinear, and simulation of the effect is indispensable. From simulation results, a scaling law for H_D has been deduced for round beams (i.e., $R = \sigma_x/\sigma_y = 1$):[11]

$$H_D = 1 + D^{1/4}\left(\frac{D^3}{1+D^3}\right)\left\{\ln(\sqrt{D}+1) + 2\ln(0.8/A)\right\} \quad , \tag{2.4}$$

where $A \equiv \sigma_z/\beta^*$, and β^* is the Courant-Snyder β-function at the interaction point. The accuracy of this scaling law is $\sim 10\%$. Thus for round beams, the effective beam size is roughly $\bar{\sigma} \sim \sigma H_D^{-1/2}$. For very flat beams (i.e., $R = \sigma_x/\sigma_y \gg 1$) and $D_x \ll 1$, however, the enhancement factor turns out to be roughly the cube-root of eq.(2.4) instead, with D and A replaced by D_y and $A_y = \sigma_z/\beta_y^*$, respectively. As is well-known, the field strength in a flat beam is largely determined by σ_x, not σ_y. So unless there is a sizable x-disruption, the mutual bootstrap of pinching between the two dimensions is lacking, resulting in a significantly milder luminosity enhancement for the flat beams.

Based on the above arguments, we deduce the following empirical rules:

$$\bar{\sigma}_x \sim \sigma_x H_{D_x}^{-1/2} \quad , \qquad \bar{\sigma}_y \sim \sigma_y H_{D_y}^{-1/3} \quad , \qquad (R \gg 1 \quad , \quad D_x \lesssim 1) \quad . \tag{2.5}$$

As can be seen from Table 1, all of the the most recent designs for the next generation linear colliders involve flat beams. Although CLIC and TESLA have $D_x \gtrsim 1$, we shall still apply eq.(2.5) as rough estimates. VLEPP has a different final focusing scheme, and our discussion above does not apply to the y-disruption for this machine. Nevertheless, its x-disruption still subject to the same condition. We emphasize that these scaling laws serve to conveniently estimate the pinch effect. For better accuracies one should resort to simulations.

Having effective beam sizes deduced, the beamstrahlung parameter is therefore

$$\Upsilon = \frac{5}{6}\frac{r_e^2\gamma N}{\alpha\sigma_z(\bar{\sigma}_x + \bar{\sigma}_y)} \quad . \tag{2.6}$$

In terms of the beamstrahlung parameter, the rate of radiating photons with energy x can be derived,

$$\nu(x) = \frac{1}{1-x}\int_x^1 dx'[x'\nu_{cl} + (1-x')\nu_\gamma] = \frac{1}{2}\Big[(1+x)\nu_{cl} + (1-x)\nu_\gamma\Big] \quad , \tag{2.7}$$

where

$$\nu_{cl} = \frac{5}{2\sqrt{3}}\frac{\alpha^2}{r_e\gamma}\Upsilon \quad , \qquad \nu_\gamma = \nu_{cl}[1+\Upsilon^{2/3}]^{-1/2} \quad . \tag{2.8}$$

With these basic parameters introduced, $f_r(x)$ is given by[12]

$$f_r(x) = \frac{1}{\Gamma(1/3)}\Big(\frac{2}{3\Upsilon}\Big)^{1/3} x^{-2/3}(1-x)^{-1/3}\exp\Big[-\frac{2x}{3\Upsilon(1-x)}\Big]\cdot G(x) \ , \tag{2.9}$$

where

$$\begin{aligned} G(x) &= \frac{1-w}{g(x)}\Big\{1 - \frac{1}{g(x)n_\gamma}\Big[1 - e^{-g(x)n_\gamma}\Big]\Big\} + w\Big\{1 - \frac{1}{n_\gamma}\Big[1 - e^{-n_\gamma}\Big]\Big\} \quad , \\ g(x) &= 1 - \frac{D}{\nu_\gamma}(1-x)^{2/3} \quad , \end{aligned} \tag{2.10}$$

and $w = (1/6)\sqrt{3\Upsilon/2}, n_\gamma = \sqrt{3}\sigma_z\nu_\gamma$; n_γ is the mean number of photons radiated per electron throughout the collision. The spectrum (2.9) applies for $\Upsilon \lesssim 5$.

3. The QED Backgrounds

Although the coherent pair production may be abundant beyond certain threshold, there is nevertheless a way to stay below this threshold by properly adjusting the beam parameters. On the other hand, the incoherent pairs with inherently large angles can not be avoided. All these issues have been studied in some details in recent years.[2,3,4] In this chapter we shall only breifly review the problem.

3.1 Coherent Pair Creation

A photon in vacuum is always accompanied with virtual electron-positron pairs. When the photon traverses a strong transverse electromagnetic field, however, the energy-momentum can be carried by the field and the pair can be kicked on-shell. Consider the boosted frame where the e^+e^- pair is created at rest. In this frame there is an electric field which is $E' = (\hbar\omega/2m_ec^2)B$, where B is the magnetic field in the lab frame. At the threshold, the created particle with unit charge e should acquire enough energy within one Compton wavelength to supply for its rest mass. Thus the threshold condition is $eE'\lambda_e \sim m_ec^2$, or $(\omega/m_e)B/B_c \sim 1$. Accordingly, there exists a minimum energy, ε_{min} in the spectrum, which, in contrast to the incoherent case, is much larger than the electron rest mass: Again in the Lorentz frame where the pair is created at rest, the invariant mass of the system is $W = 2eE'\lambda_e$. The Lorentz factor for the boost is obviously the photon energy ω devided by the invariant mass. Thus we have $W^2 = 2eB\omega\lambda_e$. On the other hand, from the final state we have $W^2 = \omega^2m_e^2/\varepsilon_+\varepsilon_-$, where $\varepsilon_+, \varepsilon_-$ are the energies of the pair particles. In the case where one particle is at very low energy, e.g., $\varepsilon_+ \ll \varepsilon_- \sim \omega$, we have $W^2 \sim \omega m_e^2/\varepsilon_+$. Thus $\varepsilon_{min} \sim \gamma m_e/2\Upsilon$. The actual value of ε_{min} is somewhat different from this naive picture and is $\sim \gamma m_e/10\Upsilon$.

The total number of coherent pairs created per primary beam particle is found to be

$$n_c = \left(\frac{\alpha\sigma_z}{\gamma\lambda_c}\Upsilon\right)^2 \Xi(\Upsilon) \quad , \tag{3.1}$$

where

$$\Xi(\Upsilon) = \begin{cases} (7/128)\exp(-16/3\Upsilon) \quad , & (\Upsilon \lesssim 1) \quad ; \\ 0.295\Upsilon^{-2/3}(\log\Upsilon - 2.488) \quad , & (\Upsilon \gg 1) \quad . \end{cases} \tag{3.2}$$

It turns out that in linear collider designs the quantity $(\alpha\sigma_z/\gamma\lambda_c)\Upsilon$ is not arbitrary. In order that the average energy loss through beamstrahlung, δ_B, is below 10 to 20 %, $(\alpha\sigma_z/\gamma\lambda_c)\Upsilon$ is constrained to be of order unity. We can thus see from the above expression that $n_b \sim \mathcal{O}(10^{-2})$ for $\Upsilon \gtrsim 1$, while for $\Upsilon \lesssim 1$, the number of pairs is exponentially suppressed. Since the typical number of particles in a bunch is $\sim \mathcal{O}(10^{10})$, we expect to have $\sim \mathcal{O}(10^8)$ e^+e^- pairs per collision in the $\Upsilon \gtrsim 1$ regime, and have the pairs totally suppressed if $\Upsilon \lesssim 0.3$.

3.2 Incoherent Pair Creation

The partial cross section for the pair-created positron with transverse momentum $p_\perp \geq p_*$ and outcoming angle $\theta_0 \leq \theta \leq \pi - \theta_0$ is

$$\sigma_{e^+e^-}(p_*, \theta_0) = \int_{-c_0}^{c_0} dc \int_{x_-}^{\infty} dx_2 \int_{x_b}^{\infty} dx_1 L_{\gamma\gamma}(x_1, x_2) \cdot \sigma(\gamma(x_1)\gamma(x_2) \rightarrow e^+e^-) \quad , \tag{3.3}$$

where $c_0 \equiv \cos\theta_0$, $x_b = x_2x_+/(x_2 - x_-)$, $x_\pm = (p_*/2\gamma m_e)\sqrt{(1 \pm c)/(1 \mp c)}$, and x_1, x_2 are the fractions of the total energy of the initial electrons and positrons, respectively, carried by the colliding photons. As noted in the introduction, the luminosity function receives contributions from two sources, bremsstrahlung and beamstrahlung, corresponding to real and virtual photons. Assuming that the sources of the two photons are independent of one another, we can write the luminosity functions as a sum of components:

$$L_{\gamma\gamma}(x_1,x_2) = f_r(x_1)f_r(x_2) + \left[f_v(x_1)f_r(x_2) + f_r(x_1)f_v(x_2)\right] + f_v(x_1)f_v(x_2). \quad (3.4)$$

In this equation, $f_v(x)$ is the Weiszacker-Williams distribution for radiation in a collision process, $f_r(x)$ is the average of the beamstrahlung spectrum over the process of interpenetration of the e^- and e^+ bunches. The three contributions in $L_{\gamma\gamma}(x_1,x_2)$ corresponds to the Breit-Wheeler, Bethe-Heitler, and Landau-Lifshitz processes, respectively. Using

$$f_v(x) = \frac{2\alpha}{\pi}\frac{1}{x}\ln\left(\frac{1}{x}\right) \quad , \quad (3.5)$$

and an approximate, single-photon limit, of the beamstrahlung spectrum

$$\lim_{n_\gamma \to 0} f_r(x) = \frac{1}{2\pi}\Gamma(2/3)\left(\frac{\alpha\sigma_z}{\gamma\lambda_c}\right)(3\Upsilon)^{2/3}y^{-2/3} \equiv Ay^{-2/3} \quad , \quad (3.6)$$

where $\Gamma(2/3) \simeq 1.3541$, and with the cross section for $\gamma\gamma \to e^+e^-$:

$$\sigma(\gamma\gamma \to e^+e^-) \approx \frac{\pi r_e^2}{\gamma^2 x_1 x_2}\frac{1}{1-c^2}, \quad (3.7)$$

it is found that[13]

$$\sigma_{e^+e^-}(p_*, \theta_0) = \sigma_{BW} + \sigma_{BH} + \sigma_{LL} \quad , \quad (3.8)$$

with

$$\sigma_{BW} = 1.69\frac{r_e^2}{\gamma^2}A^2\left(\frac{2\gamma m_e}{p_*}\right)^{4/3}\log\frac{1}{\tau_0} \quad ;$$
$$\sigma_{BH} = 3.55\frac{\alpha r_e^2}{\gamma^2}A\left(\frac{2\gamma m_e}{p_*}\right)^{5/3}\left[\tau_0^{1/3} - \tau_0^{-1/3}\right]\left[\log\frac{p_*}{2\gamma m_e} + 0.21\right] \quad ;$$
$$\sigma_{LL} = 0.83\frac{\alpha^2 r_e^2}{\gamma^2}\left(\frac{2\gamma m_e}{p_*}\right)^2\log\frac{1}{\tau_0}\left[\log\frac{p_*\tau_0}{2\gamma m_e}\log\frac{p_*}{2\gamma m_e \tau_0} + 3\log\frac{p_*}{2\gamma m_e} + 4.44\right] \quad ;$$

where $\tau_0 = \tan(\theta_0/2)$. The above expressions account for only one of the two particles (say the positron) in the pair. To count electrons as well, we must multiply each expression by 2.

It turns out that for processes involving virtual photons, the region of the impact parameter (the inverse of the transverse momentum transfer) which is larger than the beam size will be suppressed.[14] Effectively, this geometric reduction effect modifies the virtual photon spectrum into $f_v(x) \simeq (2\pi/\alpha)(1/x)\ln(2\sigma_y/\lambda_c)$. One can in principle repeat the calculations with this spectrum. The details are beyond the scope of this paper. Roughly speaking, at $p_* = 20$ MeV and $\theta_0 = 0.15$, the reduction is about 40% for very flat beams like that in NLC and JLC.[13]

4. The QCD Backgrounds

As discussed in the Introduction, photons also resolve into partons and interact hadronically. The hard scatterings between the partons will result in the form of minijets, which would be another souce of backgrounds. The cross section is again describable in the form of

$$\sigma(e^+e^- \to X + \text{anything}) = \int_0^1 dx_1 \int_0^1 dx_2 L_{\gamma\gamma}(x_1, x_2) \cdot \sigma(\gamma(x_1)\gamma(x_2) \to X) \quad . \tag{4.1}$$

To compute the jet production cross section at a jet transverse momentum of order Q, Dress and Godbole have argued that one should use a modified version of the standard Weiszacher-Williams formula. The standard formula integrates over all photon transverse momenta, as in the case of incoherent pair creation; however, only those photons which are off-shell by less than Q^2, and only a fraction of those, will contain partons which can produce jets by scattering from partons of the target. Following this argument, we take[5] in this case

$$f_v(x) = c_v \cdot \frac{\alpha}{2\pi} \frac{1+(1-x)^2}{x} \log \frac{Q^2}{m_e^2} \quad . \tag{4.2}$$

where $c_v = 0.85$. Unlike the e^+e^- pair creation process, the cross term in $L_{\gamma\gamma}$ in this case does not suffer any geometric reduction because of the typical largeness of Q.

While there is no essential disagreement on $L_{\gamma\gamma}$, the jet cross section $\sigma(\gamma\gamma \to X)$ has been a subject of debate. To elucidate the point, let us define the jet yield $\mathcal{Y}(p_*)$ as the expected number of jets with $p_\perp > p_*$, divided by the luminosity. The jet yield $\mathcal{Y}(p_*)$ can be computed from the formula

$$\mathcal{Y}(p_*) = \int_0^1 dz_1 F(z_1) \int_0^1 dz_2 F(z_2) \int_{-1}^1 dc \frac{d\sigma}{dc}(gg \to gg) \cdot \theta(p_\perp - p_*) \quad . \tag{4.3}$$

In this formula, the parton-parton scattering angle is measured in the center-of-mass frame. Let us take the parton distribution $F(z)$ to be the sum of gluon and quark distributions

$$F(z) = f_g(z) + \frac{4}{9}\sum_i (f_{qi}(z) + f_{\bar{q}i}(z)) \quad , \tag{4.4}$$

with the appropriate coefficient that we can approximate all of the parton cross sections by the gluon-gluon cross section:

$$\frac{d\sigma}{dc}(gg \to gg) = \frac{9}{16}\frac{\pi\alpha_s^2}{\hat{s}}\left[\frac{(2+\cos^2\theta)^3}{\sin^4\theta}\right], \tag{4.5}$$

where $\hat{s} = z_1 z_2 s$ is the square of the gluon-gluon center of mass energy. The coupling constant α_s is evaluated at the momentum scale $p_\perp$.

Using the Drees-Grassie parametrization[15] for the parton distributions of the photon, and with $\alpha_s(3\text{ GeV}) = 0.37$, it is found that the dependence of the jet yield on energy and p_* is well described by the parametrization[6]

$$\mathcal{Y}(p_*, E_{\text{cm}}) = A_1 \frac{(E_{\text{cm}})^{A_2}}{(A_3 + p_*)^2} \exp\left\{-\frac{B(p_*)}{(E_{\text{cm}} - p_*)^{C(p_*)}}\right\} \quad , \tag{4.6}$$

with $A_1 \approx 4000$, $A_2 = 0.82$, $A_3 = 3.0$, and

$$B(p_*) = 14.2\tanh(0.43 p_*^{1.1}) \; , \qquad C(p_*) = 0.48 p_*^{-0.45} \; . \tag{4.7}$$

E_{cm} and p_* are in units of GeV. This parametrization fits the numerical evaluation to within 20% accuracy for $p_* < 10$ GeV and $E_{\text{cm}} < 10$ TeV. We shall use this paramerization in the following discussions. With various sources of uncertainties, we expect that it yields a calculation of $\mathcal{Y}(p_*)$ up to an uncertainty of about a factor of 2.

4.1. The $\gamma\gamma$ Total Cross Section

In essence, the "minijet model" (MJ) of the total cross section would be to take

$$\sigma(\gamma\gamma \to X) = \sigma_0 + \frac{1}{2}\mathcal{Y}(p_*), \tag{4.8}$$

where σ_0 is a constant soft-scattering cross section and the cutoff p_* is taken sufficiently large that events contributing to the jet yield are not also accounted as part of σ_0. This is not exactly the model advocated by Drees and Godbole; they omit the constant term, and, at the end of ref. 6, they argue that the jet yield estimate should be modified in a manner similar to what we have described above. If it does not include the effects of soft hadronic reactions, the prediction for the cross section will be too small at low energy.

Earlier, it has been argued that the photon cross sections cannot rise as fast as the jet yield is predicted to rise.[16,17] The easiest way to argue to this conclusion is to apply this prescription for $p\bar{p}$ collisions and compare the results to the data on the $p\bar{p}$ total cross section. One finds that[9] the jet yield calculation using a value of $p_* = 1.6$ GeV, which was used by Drees and Godbole[4], is completely incompatible with the $p\bar{p}$ total cross section in a region where this cross section is well measured.

Notice that for any value of p_*, the MJ prediction for the cross section rises much faster at high energy than the expectation from the vector dominance picture. In order to produce a significantly larger cross section than this, either the photon must become larger or it must become a hadron with higher probability. Resolving the hadronic components of the photon into partons does not increase the size of the photon. Altarelli-Parisi evolution can create new hadronic components of the photon, through the diagram in which the photon off shell by an amount Q splits to a $q\bar{q}$ pair. This diagram has a substantial effect on the total number of gluons in the photon, but it has only a small effect on the photon's hadronic cross section, since the new hadronic component has the very small size π/Q^2. It is possible to explain a slowly rising cross section by making a model in which the soft hadron is a grey scattering distribution which becomes black as the gluon-gluon scattering becomes important. As the disk becomes black, the effect of gluon-gluon scattering on the total cross section must turn off. This physical effect can be implemented in a calculational scheme called 'eikonalization'. For the case of γp scattering, models of this sort have been constructed by Durand and Pi,[16] Forshaw and Storrow,[18] and Fletcher, Gaisser and Halzen.[19] Forshaw and Storrow have also written an eikonalized model of the $\gamma\gamma$ cross section.[7]

The Reference Model[9] follows the same philosophy, and takes the parametrization of Amaldi *et al.*[20] as a first approximation to the energy-dependence of the cross section for hadron production in $\gamma\gamma$ collisions:

$$\sigma_{\rm had} \equiv \sigma(\gamma\gamma \to \text{hadrons}) = \sigma_0\Big[1 + (6.30\times 10^{-3})\{\log(s)\}^{2.1} + (1.96)s^{-0.37}\Big], \quad (4.9)$$

where s is given in $(\text{GeV})^2$. The constant is adjusted so that $\sigma(\gamma\gamma) = [\sigma(\gamma p)]^2/\sigma(pp)$ in the region of approximately constant cross sections at $E_{\rm cm} \sim 30$ GeV: $\sigma_0 = 200$ nb. Comparing $\sigma(\gamma p)$ to $\sigma(\pi p)$, we conclude that the photon is a hadron a fraction (1/300) of the time.

4.2. Minijet Yields

To a first approximation, the jet yield $\mathcal{Y}(p_*)$ computed from eq.(4.3) should be a valid estimate of the total number of jets produced even when the jet yield substantially overestimates the total hadronic cross section. The reason for this is that the individual parton-parton interactions are relatively weak, and it is only because there are many gluons in a hadron that the sum of these cross sections saturates the geometrical limit on the cross section. In other words, those events in which the hadronic disks overlap typically contain a soft interaction plus gluon-gluon scatterings; if $\mathcal{Y}(p_*) \gg \sigma_{had}$, typical encounters contain many individual gluon-gluon collisions. If we assume that these collisions are completely independent, we would expect the number of pairs of of jets per event to follow a Poisson distribution, such that the mean number of jets per event is

$$\langle n_{jet} \rangle = \mathcal{Y}(p_*)/\sigma_{had} \quad . \tag{4.10}$$

The cross section for events with jets of $p_\perp > p_*$, in the Reference Model, is

$$\sigma_{jet}(p_*) = \sigma_{had} \cdot \left\{1 - \exp\left[-\mathcal{Y}(p_*)/2\sigma_{had}\right]\right\} \quad . \tag{4.11}$$

The combination of these ideas has an interesting implication. $\mathcal{Y}(p_*)$ increases much more rapidly with energy than σ_{had}. However, in this picture, the main effect of the increase in $\mathcal{Y}(p_*)$ is not to increase the hadronic cross section but rather to increase the number of jets per event. For photon-photon collisions, and for hadron-hadron collisions, above 1 TeV in the center of mass, we expect that the typical event is bristling with jets of 10 GeV transverse momentum. The time structure of jet events, in this veiw, is not evenly smeared at every e^+e^- beam collision. Instead, it bursts once in a while with high multiplicity. This casts the problem of hadronic jets underlying e^+e^- annihilation events in a quite different form, which is probably much easier to ameliorate.

5. Linear Collider Parameters

We now estimate the various QED and QCD backgrounds for the 0.5 TeV linear colliders currently under study. All designs except CLIC involve $\Upsilon < 0.3$, and the coherent pairs are totally suppressed. CLIC would yield a total of $N_c = 413$ coherent pairs per bunch crossing. The number of incoherent pairs per bunch crossing, $N_{e^+e^-}$, is calculated using eq.(3.8) with $p_* = 20$ MeV and $\theta_0 = 0.15$. The geometric reduction is not included. At this choice of angular-momentum cuts, the reduction is about 40% for the smallest beam sizes like in NLC, JLC, and VLEEP, and milder for other machines. For the minijet events per bunch crossing, N_{jet}, we take $p_* = 3.2$ GeV and 8 GeV. It was shown[8] that a choice of $p_* = 3.2$ GeV fits the UA1 minijet data at 5 GeV transverse energy. So we interpret the calculated N_{jet} at $p_* = 3.2$ and 8 GeV as that for 5 and 10 GeV transverse energies.

From the Table we see that for e^+e^- colliders at 0.5 TeV, neither e^+e^- nor minijet backgrounds look severe. However, for these machines and certainly for future colliders, it is important to learn what parameters of the $\gamma\gamma$ event spectrum do constrain the experimental environment and must be minimized in any design. It seems likely that only those events of sufficiently large $\gamma\gamma$ collision energy or jet transverse momentum will be a serious problem[21].

Table 1. Parameters and Backgronds for 0.5 TeV Linear Colliders

Linear Colliders	CLIC	DLC	JLC	NLC	TESLA	VLEPP
$L_0[10^{33}\mathrm{cm}^{-2}\mathrm{sec}^{-1}]$	2.7	2.4	6.8	6.0	2.6	12
$f_{rep}[\mathrm{Hz}]$	1700	50	150	180	10	300
n_b	4	172	90	90	800	1
$L_0/(f_{rep}\cdot n_b)[10^{30}\mathrm{cm}^{-2}]$	0.40	0.27	0.50	0.37	0.33	40
$N[10^{10}]$	0.6	2.1	0.7	0.65	5.15	20
$\sigma_x/\sigma_y[\mathrm{nm}]$	90/8	400/32	260/3	300/3	640/100	2000/4
$\sigma_z[\mu m]$	170	500	80	100	1000	750
$\beta_x^*/\beta_y^*[\mathrm{mm}]$	2.2/0.16	16/1	10/0.1	10/0.1	10/5	100/0.1
D_x/D_y	1.3/15	0.70/8.7	0.07/6	0.08/8.2	1.25/8.0	0.43/—
A_x/A_y	0.08/1.06	0.03/0.5	0.008/0.8	0.01/1.0	0.1/0.2	0.008/—
$\bar{\sigma}_x/\bar{\sigma}_y[\mathrm{nm}]$	40/5.5	246/19	259/2.0	300/2.2	304/50	1587/4
H_D	3.3	2.8	1.5	1.4	4.2	1.26
$L[10^{33}\mathrm{cm}^{-2}\mathrm{sec}^{-1}]$	8.85	6.55	10.0	8.2	11.1	15.1
Υ_0	0.16	0.043	0.15	0.095	0.031	0.059
Υ	0.35	0.071	0.15	0.096	0.065	0.076
δ_B	0.36	0.08	0.05	0.03	0.14	0.14
n_γ	4.6	3.1	1.0	0.84	5.8	5.1
$N_{e^+e^-}(p_t = 20\mathrm{MeV})$	23.4	14.0	4.8	3.2	54.6	1564
N_{had}	1.35	0.29	0.06	0.03	1.53	45.5
$N_{\mathrm{jet}}(p_t = 5\mathrm{GeV})[10^{-2}]$	5.97	0.43	0.22	0.10	1.61	5.83
$N_{\mathrm{jet}}(p_t = 10\mathrm{GeV})[10^{-4}]$	17.06	1.14	0.68	0.31	3.89	114.8

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