PROMPT PHOTINO PRODUCTION AT THE CERN $\bar{p}p$ COLLIDER ENERGIES AND BEYOND

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We present a lowest-order calculation of the inclusive cross section for a prompt photino to be produced in association with either an ordinary QCD jet or a photon in $\bar{p}p$ collisions. If the gluino and squark masses are not too large ($m \lesssim E_T^{\text{trigger}}$), so that the supersymmetric content of the proton – controlling the subprocesses considered – is plausibly excited, we find missing-energy rates detectable in the present data from the CERN $\bar{p}p$ collider. These rates increase dramatically in the TeV energy range.

Apart from its mathematical elegance, supersymmetry (SUSY) [1] may provide solutions to fundamental theoretical problems, such as the hierarchy of the symmetry-breaking scales, and the eventual incorporation of gravity. The most direct phenomenological consequence of SUSY is the existence of a bosonic (fermionic) partner to every known fermion (boson) [2]. Although the elementary vertices involving these superpartners are completely determined by the supersymmetry of the lagrangian, the masses they acquire after the strong SUSY breaking required by low-energy ($\lesssim$1 TeV) phenomenology are very much model dependent. Despite this uncertainty, interesting phenomenological implications of wide classes of SUSY theories may be formulated by treating the masses of the superpartners as free parameters, on which the data may provide information [3].

A general experimental signature of SUSY at presently accessible energies is a large apparent missing energy in the final state. Indeed, in most SUSY theories the lightest sparticle, in which all heavier superpartners decay, is the photino, which is not normally detected since it is neutral, not strongly interacting and stable (very much like the neutrino). Thus, the unusual events recently observed at the CERN $\bar{p}p$ collider [4] to be marked with characteristic missing energy, if proved statistically significant, may constitute important evidence for sparticle production [4].

Since gluino production from the subprocesses gg, $\bar{q}q \to \tilde{g}\tilde{g}$ is significant, the decay $\tilde{g} \to q\bar{q}\tilde{\gamma}$, though of higher order in α_s, is found to give rise to observable signatures in the present collider data for sufficiently light gluinos, $m_{\tilde{g}} \lesssim 50$ GeV [5]. On top of that, in analogy to direct photons of ordinary QCD, direct photinos may also be produced in $\bar{p}p$ collisions. In this work we make a clear lowest-order theoretical calculation of the inclusive photino production cross section, where a direct photino is produced in association with either an ordinary QCD quark or gluon jet or a photon. The relevant subprocesses, depicted in fig. 1, are

$$\tilde{g}q \to \tilde{\gamma}q\,, \tag{1a}$$

$$\tilde{q}g \to \tilde{\gamma}q\,, \tag{1b}$$

$$\tilde{q}\bar{q} \to \tilde{\gamma}g\,, \tag{1c}$$

$$\tilde{q}\bar{q} \to \tilde{\gamma}\gamma\,. \tag{1d}$$

They necessarily involve either a gluino or a squark in the initial state. We assume that these emerge from the proton's SUSY content, which is plausibly excited at high values ($\gtrsim 4m^2$) of the probing large variable Q^2, i.e. at sufficiently large transverse energies ($E_T^2 \sim Q^2$) [6,7]. Although the gluino and the squark distributions in the proton are small, for $m_{\tilde{q}}$, $m_{\tilde{g}} \sim 30$ GeV we find a $\tilde{\gamma}$-production rate (pro-

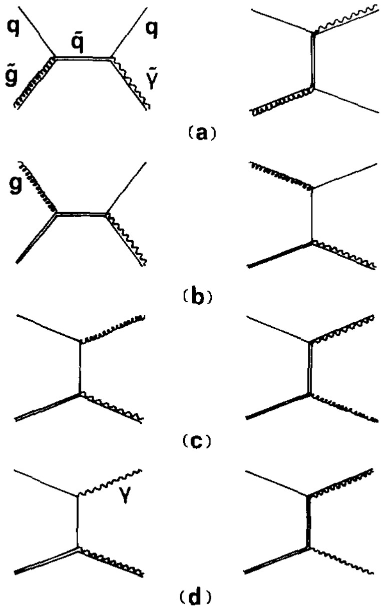

Fig. 1. (a)–(c): Lowest order SUSY–QCD graphs leading to photino plus ordinary QCD jet in the final state. (d): Lowest order SUSY–QED graph resulting in photino plus photon production.

viding apparent missing energy in the final state) detectable in the present UA1 data [4]. We predict that this rate will increase by about one order of magnitude at $E_{CM} \sim 1$ TeV and by more than two orders at $E_{CM} \sim 10$ TeV, for $E_T^{\tilde{\gamma}} \sim 50$ GeV. We next briefly describe the calculation:

Using the relevant lagrangians [2,6] it is straightforward to calculate the graphs of fig. 1. The cross sections (1a)–(1d), with a massless photino and mass degenerate left- and right-handed squarks, are

$$\frac{d\hat{\sigma}}{dt}(\tilde{g}q \to \tilde{\gamma}q) = \frac{\pi\alpha\alpha_s}{(s - m_{\tilde{g}}^2)^2} 32 \times \left(\frac{s(s - m_{\tilde{g}}^2)}{(s - m_{\tilde{q}}^2)^2} + \frac{u(u - m_{\tilde{g}}^2)}{(u - m_{\tilde{q}}^2)^2}\right), \tag{2a}$$

$$\frac{d\hat{\sigma}}{dt}(\tilde{q}g \to \tilde{\gamma}q) = \frac{\pi\alpha\alpha_s}{(s - m_{\tilde{q}}^2)^2} 32 \times \left(-\frac{s(s + m_{\tilde{q}}^2)}{(s - m_{\tilde{q}}^2)^2} + \frac{s(t - 2m_{\tilde{q}}^2)}{(s - m_{\tilde{q}}^2)t} + \frac{tu}{t^2}\right), \tag{2b}$$

$$\frac{d\hat{\sigma}}{dt}(\tilde{q}\bar{q} \to \tilde{\gamma}g) = \frac{\pi\alpha\alpha_s}{(s - m_{\tilde{q}}^2)^2} 32 \times \left(-\frac{st}{t^2} + \frac{u(t - 2m_{\tilde{q}}^2)}{t(u - m_{\tilde{q}}^2)} + \frac{u(u + m_{\tilde{q}}^2)}{(u - m_{\tilde{q}}^2)^2}\right), \tag{2c}$$

$$\frac{d\hat{\sigma}}{dt}(\tilde{q}\bar{q} \to \tilde{\gamma}\gamma) = \frac{\pi\alpha^2}{(s - m_{\tilde{q}}^2)^2} 24 \times \left(-\frac{st}{t^2} + \frac{u(t - 2m_{\tilde{q}}^2)}{t(u - m_{\tilde{q}}^2)} + \frac{u(u + m_{\tilde{q}}^2)}{(u - m_{\tilde{q}}^2)^2}\right), \tag{2d}$$

where s, t, u are the subprocess Mandelstam invariants: $s + t + u = m_{\tilde{g}}^2$ in (2a) or $= m_{\tilde{q}}^2$ in (2b)–(2d). Note that in the first graph of fig. 1a, in which a physical-region pole appears for $m_{\tilde{g}} < m_{\tilde{q}}$, the squark lifetime becomes relevant: we take $\Gamma_{\tilde{q}} \sim (m_{\tilde{q}}$ in GeV) $\times 10^{-21}$ s^{-1} [8] but our results are insensitive to this for $m_{\tilde{q}} \lesssim 60$ GeV and $E_T^{\tilde{\gamma}} \gtrsim 35$ GeV.

The coupling strength of supersymmetric QCD (SQCD) is given by [6]

$$\alpha_s = [12\pi/(33 - 2f - 2N_\lambda)]\,(\ln Q^2/\Lambda_1^2)^{-1}, \quad \text{for } 2m_{\tilde{g}} \leqslant Q \leqslant 2m_{\tilde{q}}, \tag{3a}$$

$$= [12\pi/(33 - 2f - f)]\,(\ln Q^2/\Lambda_1^2)^{-1}, \quad \text{for } 2m_{\tilde{q}} \leqslant Q \leqslant 2m_{\tilde{g}}, \tag{3b}$$

$$= [12\pi/(33 - 2f - 2N_\lambda - f)]\,(\ln Q^2/\Lambda_2^2)^{-1}, \quad \text{for } Q \leqslant m_{\tilde{g}}, m_{\tilde{q}}, \tag{4}$$

where f is the number of flavours and $N_\lambda = 3$ for Majorana gluinos. The effective scale parameters Λ_1, Λ_2 are determined by matching (3) to the coupling strength of the ordinary QCD and to (4) at the appropriate thresholds. We take $\Lambda_{QCD} = 0.2$ GeV; our results do not depend crucially on this.

To calculate the contribution of the subprocesses (2) to the observable jet cross section we fold them with parton distributions according to the factorizable formula

$$\left[\frac{d\sigma^{\tilde{\gamma}}}{dE_T^{\tilde{\gamma}}\, d\eta^{\tilde{\gamma}}}\right]_{\eta^{\tilde{\gamma}}=0} = 2E_T^{\tilde{\gamma}} \sum_{i,j} \int_{x_i^{\min}}^{1} \frac{(\rho + r)\, dx_i}{(\rho + r)x_i - x_T}$$

$$\times\, G_i(x_i) G_j(x_j) \frac{d\hat{\sigma}}{dt}(s, t)\,, \tag{5}$$

where

$$r = (1 - \epsilon)^{1/2}, \quad \rho = [1 + \epsilon(1 - x_i^2)/x_i^2]^{1/2}\,, \tag{6}$$

$$\epsilon = 4m^2/E_{CM}^2, \quad x_T = 2E_T^{\tilde{\gamma}}/E_{CM}\,, \tag{7}$$

$$s = m^2 + \tfrac{1}{2}(\rho + r)x_i x_j E_{CM}^2\,, \tag{8}$$

and $m = m_{\tilde{g}}$ for (2a) or $m = m_{\tilde{q}}$ for (2b)–(2d). At $\theta_{CM} = 90°$ ($\eta^{\tilde{\gamma}} = 0$) we have

$$t = m^2 - \tfrac{1}{2}\rho x_i x_T E_{CM}^2\,, \tag{9}$$

$$x_j = \frac{-\epsilon/2 + \rho x_i x_T}{(\rho + r)x_i - x_T} \tag{10}$$

$$x_i^{\min} = \frac{x_T - \epsilon/2 - (1 - x_T)[(x_T + \epsilon/2) - \epsilon x_T^2]^{1/2}}{x_T(2 - x_T)(1 - \epsilon)^{1/2}}\,. \tag{11}$$

To proceed we require the SUSY sea distributions in the nucleon. We take

$$x\tilde{g}(x) = [4(1 + \lambda)/(5f + 20)](1 - x)^{\lambda}\,, \tag{12}$$

which gives the correct total momentum carried by the gluinos at infinite Q^2 [6]. We take $\lambda \simeq 25$, as suggested by the results of refs. [6]. This is in variance with ref. [7], where a much harder gluino distribution is considered. We also used a Q^2-dependent gluino distribution [9], which was obtained by iterating the Altarelli–Parisi equations, appropriately generalized for SQCD with gluinos [6]: The latter leads to a flatter cross section but both give comparable results (within a factor of 2) for $E_T^{\tilde{\gamma}} \sim 30$–50 GeV.

Similarly, neglecting the valence squark components and not discriminating between different squark flavours, we take

$$x\tilde{q}'(x) = [(1 + \mu)/(5f + 20)](1 - x)^{\mu}\,, \tag{13}$$

with $\mu \simeq 25$ as again suggested by the results of refs. [6]. Taking $\lambda, \mu = 20$ (30) results to approximately double (half) cross sections at $E_T^{\tilde{\gamma}} \sim 60$ GeV. It should be noted that although the valence-squark component we neglect falls off slower than the singlet component with x, its absolute magnitude at the small-x region, which gives most of the contribution to the integral (5), is very small.

For the ordinary parton distributions we used the parametric forms of both refs. [10,11]: given the ambiguity of the SUSY sea distributions, both these sets give similar results. As "large variable" we take $Q^2 = 2(E_T^{\tilde{\gamma}})^2$: As is well known, other popular choices give results different by a factor, which can be as large as 1.5.

In this calculation we neglect the (mass) threshold effects due to the excitation of the heavy spartons $\tilde{g}$, $\tilde{q}$. We apply the *factorizable* formula (5) using the *asymptotic* ($Q^2 \to \infty$) sparton distributions (12), (13): our results are strictly valid for $E_T^{\tilde{\gamma}} \gg m$. Optimistically, we draw curves for $E_T^{\tilde{\gamma}} \gtrsim \sqrt{2}m$ taking literally the rule $Q^2 = 2(E_T^{\tilde{\gamma}})^2 > 4m^2$, but this somehow overestimates the small-$E_T^{\tilde{\gamma}}$ cross sections.

Our results with $\lambda = \mu = 25$ and with the quark–gluon distribution set of ref. [11] are summarized in figs. 2–4. Given the present UA1 integrated luminosity (~113/nb), and the experimental E_T resolution (~5 GeV), the one-event level for the present data is at ~1.8 pb/GeV. As seen from figs. 2–4 we

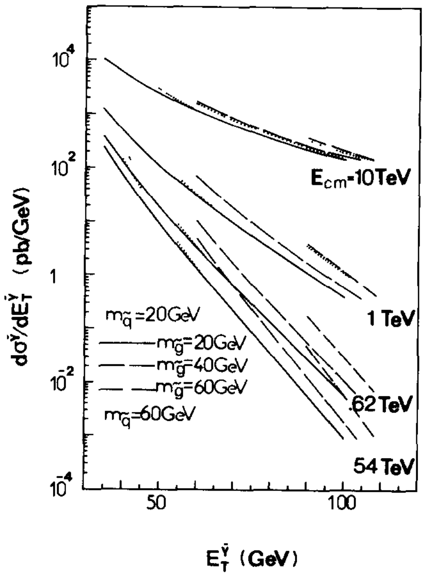

Fig. 2. Contribution of the subprocess (1a) to the inclusive photino cross section for various CM-energies and gluino–squark masses, as a function of the transverse energy of the produced photino, in $p\bar{p}$ collisions.

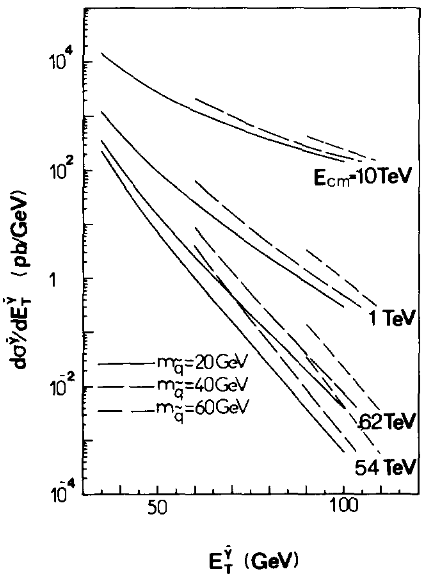

Fig. 3. Same as in fig. 2, for the sum of the subprocesses (1b) and (1c). In our model these results depend only very slightly on the gluino mass.

require sparton masses $m_{\tilde{g}}, m_{\tilde{q}} \gtrsim 30$ GeV in order that the subprocesses (1) do not give rise to too many monojet or monophoton events with small $E_T^{\tilde{\gamma}}$ at $E_{CM} = 0.54$ TeV. For $E_T^{\tilde{\gamma}} \sim 40$–$60$ GeV our monojet as well as our monophoton rate increases by about a factor of 2 when the CM-energy is increased to 0.62 TeV. The missing-energy rate $d\sigma^{\tilde{\gamma}}/dE_T^{\tilde{\gamma}}$ falls-off exponentially, with a slope $\sim 0.2\ \mathrm{GeV}^{-1}$ at $E_{CM} = 0.54$ TeV, giving a total of about $\sim \frac{1}{2}\,[d\sigma^{\tilde{\gamma}}/dE_T^{\tilde{\gamma}}$ (in pb/GeV)$]_{E_{T\min}^{\tilde{\gamma}}}$ events with $E_T^{\tilde{\gamma}} > E_{T\min}^{\tilde{\gamma}}$. One may play with our main unknowns, which are the SUSY sea exponents λ and μ, so as to get bigger or smaller cross sections. In particular, as it is clear, the ratio $\lambda : \mu$ controls the ratio of monojet : monophoton events. The poor statistics of the presently available UA1 data do not allow us to extract such specific information Our crucial conclusion is that the lowest-order processes of fig. 1, despite the smallness of the SUSY sea and independently of its details, can provide enough missing-energy cross section to be detectable at present data, provided that the spartons are light enough to be excited at present transverse-energy values.

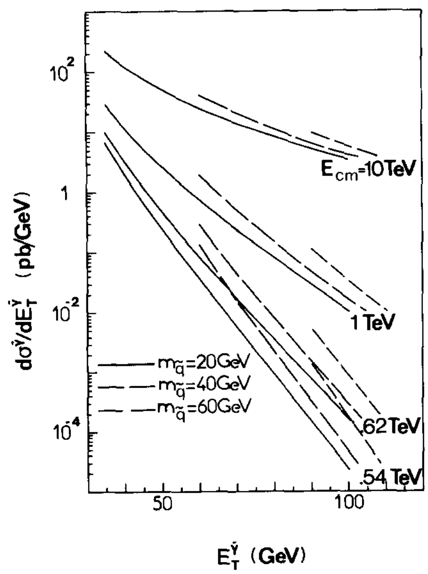

Fig. 4. Same as in fig. 3, for the subprocess (1d).

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