SEARCHING FOR A SCALAR TOP AT THE CERN $p\bar{p}$ COLLIDER

G. Altarelli

CERN, Geneva, Switzerland

Dipartimento di fisica, Univ. di Roma "La Sapienza" and
INFN, Sezione di Roma, Italy

R. Rückl

CERN, Geneva, Switzerland

ABSTRACT

In an appealing class of models of broken supersymmetry the possibility exists that the lightest top-flavoured particle is a scalar. We analyse the production and decay properties of such a top scalar especially in view of the current searches for top flavour at the CERN $p\bar{p}$ Collider.

The main task for experiment after the recent discovery of the charged and neutral gauge bosons of weak interactions is the completion of the standard model by finding the top quark and the Higgs boson(s), and the search for possible new phenomena beyond the standard theory. In this respect it is of particular importance to look for experimental clues that broken supersymmetry [1] (SUSY) plays a role in particle physics at accessible energies. On the theoretical side, considerable effort has been devoted to an appealing class of models where SUSY breaking is induced by gravitational interactions [2]. In these models the possibility exists that one of the two scalars associated with the top quark is in fact lighter than the top quark itself. Actually it could be the lightest of all scalar quarks and leptons. This is particularly likely in models where the $SU(2) \otimes U(1)$ breaking is realized at the loop-level, which requires a heavy top quark [3]. In this paper we study this possibility in detail by analysing the production and decay properties of a relatively light scalar top, especially in view of the current experiments at the CERN $p\bar{p}$ Collider where the search for the top quark is actively being carried out.

Lacking precise theoretical predictions on the pattern of masses to be expected for SUSY partners [4–6] of ordinary particles, we shall work in a schematic framework which in a simplest form takes into account all qualitative mass correlations which are in general expected in the above-mentioned class of models. In this picture of the spectrum, apart from the photino ($\tilde{\gamma}$), which we take as massless, all other masses are given in terms of the gravitino and the gluino ($\tilde{g}$) masses denoted by $m_{3/2}$ and $m_{\tilde{g}}$, respectively.

For each quark or lepton f of mass m_f we expect two associated scalar partners with masses which we assume are given by

$$m_{s_f} = | m_{3/2} \pm m_f | \quad (1)$$

With this ansatz $m_{3/2}$ must necessarily be heavier than about 20 GeV in order to accommodate the non-observation of charged scalars at PETRA and PEP. As a consequence, for all scalars associated with light fermions (including also the b quark as a first approximation) the resulting masses are almost degenerate and approximately equal to the gravitino mass:

$$m_{s_f} (f \text{ light}) \approx m_{3/2} \quad (2)$$

On the other hand, if the top quark mass m_t is comparable to $m_{3/2}$, then one of the two associated scalars is pushed down to a mass lower than both $m_{3/2}$ and m_t . This is precisely the case we want to study here.

Furthermore, if the gluino mass is small in comparison with $m_{3/2}$, we expect one of the two w -inos ($\tilde{w}$) associated with the W to be lighter [4–6] than the W itself, and similarly one of the two z -inos ($\tilde{z}$) to be lighter than the Z^0 . For the weak gaugino masses we take

$$m_{\tilde{w}, \tilde{z}} = \sqrt{m_{W, Z}^2 + \frac{1}{4} m_{3/2}^2} \pm \frac{1}{2} m_{3/2} \quad (3)$$

Equivalently, for the light w-ino and z-ino we have

$$m_{3/2} = \frac{m_W^2 - m_{\tilde{W}}^2}{m_{\tilde{W}}} = \frac{m_Z^2 - m_{\tilde{Z}}^2}{m_{\tilde{Z}}} \quad (4)$$

Note that with these mass formulae the light $\tilde{w}$ would reach 20 GeV of mass when $m_{3/2}$ is increased up to about 300 GeV. Thus the allowed range for $m_{3/2}$ is restricted to be $20 \text{ GeV} \lesssim m_{3/2} \lesssim 300 \text{ GeV}$.

In fig. 1 we report the resulting spectrum of SUSY particle masses in terms of $m_{3/2}$ for the representative case of a scalar top mass m_{s_t} of 40 GeV, a value we shall often use in the following for indicative purposes. In addition, we take $m_{\tilde{\gamma}} \approx 0$ and consider the gluino mass as a free parameter with $m_{\tilde{g}} \ll m_{3/2}$. For each given $m_{3/2}$ and a fixed scalar top mass m_{s_t} there are two possible values of the top quark mass m_t . In our discussion we always assume that $m_t > m_{s_t}$ as is, or at least may be, true for all allowed values of $m_{3/2}$. We further consider $m_{3/2} > m_{s_t}$, which is true in the case of fig. 1 in most of the range of $m_{3/2}$ ($m_{3/2} \gtrsim 40 \text{ GeV}$). Then the scalar top is the lightest of all scalar partners of the fermions. Note that for $m_{s_t} \sim 40 \text{ GeV}$ the upper limit stated for $m_{3/2}$ is also about sufficient to keep the radiative corrections to the parameter $\rho = (m_{\tilde{W}}^2/m_Z^2 \cos^2 \theta_w)$ (which determines the ratio of neutral to charged weak current rates) within a tolerable amount. These radiative corrections [7,8] are quadratically increasing with the mass differences of the top-bottom and of the scalar top-scalar bottom doublets, which therefore cannot be too large.

We immediately observe one important consequence of a general nature: whilst a top quark of 40 GeV or so would be present in the decay products of the W, because the channel $W \rightarrow t\bar{b}$ is open, owing to the large s_b mass this is not the case for a scalar top of 40 GeV, since the decay mode $W \rightarrow s_t\bar{s}_b$ is kinematically forbidden (or much suppressed in the marginal case $m_{3/2} \lesssim m_{s_t}$). For $m_{3/2} > 40 \text{ GeV}$ this actually holds for W decays into any pair of scalars. Thus, if a top-like object is detected in W decays with appreciable partial width, it has necessarily to be identified with the top quark. The Z^0 case is obviously different, and a pair of top quarks or of top scalars of sufficiently low mass would similarly contribute to Z^0 decays.

We now compare the expected production rates at the Collider for a top scalar pair and a top quark pair of the same mass. We argue that the two cross-sections are of the same order, with the quark cross-section presumably somewhat larger than the scalar quark cross-section. There is at present no reliable quantitative way of computing the total cross-sections for heavy-flavour production in hadron-hadron collisions [9]. Therefore some of our arguments will necessarily be of merely indicative nature. At present the only reliable calculation that can be made is for the production rate of heavy flavours at large transverse momentum. In this case gluon-gluon fusion and light quark-antiquark annihilation are the relevant parton subprocesses. The corresponding results for a pair of heavy quarks are well known [9]. For a pair of heavy scalars we obtained the following parton cross-sections:

$$\frac{d\sigma}{dt} (q\bar{q} \rightarrow s_t \bar{s}_t) = \frac{4\pi\alpha_s^2}{9s^4} (tu - m_{s_t}^4) \quad (5)$$

and

$$\begin{aligned}
\frac{d\sigma}{dt} (gg \rightarrow s_t \bar{s}_t) = & \frac{\pi \alpha_s^2}{4 s^2} \left\{ \frac{1}{3} \left(\frac{m_{s_t}^2 + t}{m_{s_t}^2 - t} \right)^2 - \right. \\
& - \frac{3}{16} \frac{1}{s^2} \left[2s(s - 4m_{s_t}^2) - (t-u)^2 \right] - \frac{1}{96} \frac{(s - 4m_{s_t}^2)^2}{(m_{s_t}^2 - t)(m_{s_t}^2 - u)} + \\
& + \frac{3}{32} \frac{1}{s(m_{s_t}^2 - t)} \left[(t-u)(4m_{s_t}^2 + 4t - s) - 2(m_{s_t}^2 - u) \right. \\
& \cdot (6m_{s_t}^2 + 2t - s) \left. \right] + \frac{7}{96} \frac{4m_{s_t}^2 + 4t - s}{m_{s_t}^2 - t} + \frac{7}{24} + (t \leftrightarrow u) \left. \right\}
\end{aligned}$$

where s , t , and u are the Mandelstam variables of the parton subprocesses. These results agree with those in ref. 10, but eq. (6) was never written down in full previously.

The top quark and top scalar differential cross-sections at large p_T are compared in fig. 2 for $m_t = m_{s_t} = 40$ GeV. At relatively small p_T the scalar top cross-section is suppressed with respect to the top quark cross-section, partly because the $q\bar{q}$ annihilation contribution vanishes in the forward direction for the scalar top (because of helicity conservation) whilst it stays large for the top quark. Also, both cross-sections fall steeply at large rapidities. On the other hand, we know from the comparison of this model with the data on charm production [9] that the bulk of the heavy-flavour production cross-section apparently arises from small p_T and extends to relatively large y . As already stated, the region at small p_T cannot be attacked by straightforward perturbation theory, and one has to adopt some QCD-inspired model. One possibility is to assume a given small component of heavy flavours in the nucleon, and scatter a light quark or a gluon from the other hadron off the heavy flavour in the nucleon. In these processes a massless parton is exchanged which becomes very soft as p_T decreases. Some minimum p_T must be introduced to prevent the cross-section from blowing up at $p_T = 0$. As a consequence the absolute value of the cross-section cannot be reliably predicted, because it would crucially depend on the value of the assumed minimum p_T . It is, however, reasonable to assume that the ratio of the top quark versus the top scalar cross-sections is less sensitive to this cut-off in p_T . In this spirit we compute the cross-sections for top scalar production at small but finite p_T in this model (the corresponding results for the top quark being known [11]). We obtain

$$\frac{d\sigma}{dt} (qs_t \rightarrow qs_t) = \frac{8\pi\alpha_s^2}{9t^2(s-m_{s_t}^2)^2} (m_{s_t}^2 - su) \quad (7)$$

and

$$\begin{aligned}
 \frac{d\sigma}{dt}(g s_t \rightarrow g s_t) = & \frac{\pi \alpha_s^2}{4(s-m_{s_t}^2)^2} \left\{ \frac{16}{9} \left(\frac{s+m_{s_t}^2}{s-m_{s_t}^2} \right)^2 + \right. \\
 & + \frac{1}{t^2} \left[(s-u)^2 + 2t(4m_{s_t}^2 - t) \right] + \frac{1}{18} \frac{(4m_{s_t}^2 - t)^2}{(s-m_{s_t}^2)(m_{s_t}^2 - u)} - \\
 & - \frac{1}{2} \frac{1}{(s-m_{s_t}^2)t} \left[(s-u)(4m_{s_t}^2 + 4s - t) - 2(m_{s_t}^2 - u) \right] \cdot \\
 & \cdot (6m_{s_t}^2 + 2s - u) - \frac{7}{18} \frac{4m_{s_t}^2 + 4s - t}{s-m_{s_t}^2} + \frac{14}{9} + (s \leftrightarrow u) \left. \right\} \quad (8)
 \end{aligned}$$

For $m_t \simeq m_{s_t} \simeq 40$ GeV and $p_T \lesssim 10$ GeV, the cross-sections for top quark and top scalar production turn out to be almost identical if an equal top and scalar top content in the nucleon is assumed. This near equality is a reflection of the identical vector charge form factor at vanishing momentum transfer.

On the other hand, the top quark and top scalar densities in the nucleon clearly need not be the same. However, one may reasonably assume that the ratio of the two densities can actually be determined from the ratio of the gluon fusion cross-sections. In fact the heavy-flavour content of a hadron can be imagined to arise from gluon radiation and successive materialization. From fig. 2 one reads a factor of about three between the top quark and the top scalar gluon fusion cross-sections (in the region of small y and p_T which contributes most of the cross-section). Thus we are led to conclude that the top quark and the top scalar cross-sections are expected to be of the same order, with the top quark production rate presumably larger by a factor of a few units with respect to the top scalar cross-section (at equal m_t and m_{s_t}). It also appears reasonable to assume that top quarks or scalars are produced at large y in the diffractive region with a rather small average p_T , quite negligible with respect to m_t or m_{s_t} .

We now analyse the scalar top decay modes in the proposed scenario. The diagrams for scalar top decay are depicted in fig. 3, where the intermediate particles can in some cases become real. The final state consists of a quark (presumably a b quark), a quark-antiquark or a lepton-antilepton pair, plus either a gluino or a photino:

$$\begin{aligned}
 s_t \rightarrow & b + q + \bar{q} + (\tilde{g}, \tilde{\gamma}) \\
 \rightarrow & b + \ell + \bar{\nu}_\ell + (\tilde{g}, \tilde{\gamma}) \quad (9)
 \end{aligned}$$

Note that in the semileptonic channel there is either an additional hadron jet ($\tilde{g}$) or an additional neutrino-like object ($\tilde{\gamma}$) in comparison with the analogous top quark case. Thus as a general trend the lepton spectrum is expected to be somewhat softer for a scalar top than for a top quark of equal

mass. Of special importance from the point of view of detection is the value of the branching ratio of a scalar top into a direct electron (or a muon), which we denote by $B_{SL}(s_t)$. While $B_{SL}(t)$ (for the top quark) is fixed to about 11%, $B_{SL}(s_t)$ depends on the gluino mass and on $m_{3/2}$, as we shall discuss below. However, it is easy to realize that $B_{SL}(s_t)$ can never exceed $B_{SL}(t)$, unless the assumed mass spectrum is badly violated (for example, by scalar leptons much lighter than scalar quarks).

Scalar top decays become particularly simple if the $\tilde{w}$ mass is sufficiently small such that the decay into a real $\tilde{w}$ is possible:

$$s_t \rightarrow b + \tilde{w} \quad (10)$$

The width is given by

$$\Gamma_{s_t} = \Gamma(s_t \rightarrow b \tilde{w}) = \frac{\alpha_w c_\phi^2}{4} m_{s_t} \left(1 - \frac{m_b^2}{m_{s_t}^2} - \frac{m_{\tilde{w}}^2}{m_{s_t}^2} \right) \cdot \lambda \left(1, \frac{m_b^2}{m_{s_t}^2}, \frac{m_{\tilde{w}}^2}{m_{s_t}^2} \right) \quad (11)$$

with the phase-space factor

$$\lambda(x, y, z) = \left(x^2 + y^2 + z^2 - 2xy - 2xz - 2yz \right)^{1/2} \quad (12)$$

Here, $\alpha_w = \alpha/\sin^2 \theta_w$ and c_ϕ is the projection coefficient of the physical $\tilde{w}$ on the gaugino, whilst the higgs-ino component is essentially decoupled. In the mass scheme that we have assumed, we have

$$c_\phi^2 = \frac{m_w^2}{m_w^2 + m_{\tilde{w}}^2} \quad (13)$$

For $m_{s_t} \sim 40$ GeV this decay mode is by far the dominant one in the range $m_{3/2} \geq 150$ GeV, which corresponds to $m_{\tilde{g}} \lesssim 35$ GeV. In this case the apparent $B_{SL}(s_t)$ coincides with $B_L(\tilde{w})$, the leptonic branching ratio of the $\tilde{w}$. The diagrams for $\tilde{w}$ decay can also be induced from fig. 3, by restricting ourselves to an intermediate $\tilde{w}$ on its mass shell. Note that a charged lepton from $\tilde{w}$ decay is always associated with a $\tilde{\gamma}$ and not with a $\tilde{g}$. The branching ratios of a light $\tilde{w}$ have already been studied in the literature [5,12,13]. If $\tilde{g}$ and $\tilde{\gamma}$ were degenerate in mass, the $\tilde{g}$ modes (always non-leptonic) would dominate owing to their strong coupling, and $B_L(\tilde{w})$ would be very small. However, for a light $\tilde{w}$, a moderate $\tilde{g}$ mass is sufficient to make $B_L(\tilde{w})$ comparable to $B_{SL}(t)$. For example [13], for $m_{\tilde{g}} = 10$ GeV, $B_L(\tilde{w}) \simeq (6-11)\%$ for all $m_{\tilde{g}} \lesssim 35$ GeV.

Focusing on diffractive production of scalar top, one may neglect the p_T of s_t . In that case the charged lepton p_T distribution arising from $s_t \rightarrow b\tilde{w}^+ \rightarrow b\ell^+\nu\tilde{\gamma}$ is independent of the longitudinal momentum distribution of s_t and can be directly evaluated from the relevant diagrams in fig. 3. Since the $\tilde{w}$ helicity is not observed, the charged lepton distribution cannot be obtained by an incoherent sum of convolutions of s_t decay distributions into a real $\tilde{w}$ times the lepton p_T spectrum arising from $\tilde{w}$ decay, but the absolute square of a coherent sum of helicity amplitudes has to be taken. In practice, one simply computes the relevant diagrams in fig. 3 as if the $\tilde{w}$ was virtual and then integrates over the $\tilde{w}$ Breit-Wigner, with all numerator factors evaluated on the $\tilde{w}$ mass shell. Actually, for $m_{3/2} \gtrsim 150$ GeV one can restrict the calculation to the dominant diagram of $\tilde{w}$ decay through W exchange. Neglecting the b quark mass for simplicity, the final result can be written in the form,

$$\frac{dN}{dx_T} = \frac{\alpha_w^2 \alpha C_F^4}{32\pi} \cdot \frac{m_{s_t}}{\Gamma_{s_t}} \cdot \frac{m_{\tilde{w}}}{\Gamma_{\tilde{w}}} \cdot \int_0^{\sqrt{1-x_T^2}} \frac{dx_L}{x^2} \int_x^{\min(1, x \frac{m_{s_t}^2}{m_{\tilde{w}}^2})} dy \cdot x_T \left[a + b \frac{m_{\tilde{w}}^2}{\tilde{w} \cdot \ell} \ln \frac{m_W^2}{m_W^2 - 2\tilde{w} \cdot \ell} + c \frac{m_{s_t}^4}{m_W^2 (m_W^2 - 2\tilde{w} \cdot \ell)} \right] \quad (14)$$

with

$$\begin{aligned} a &= \frac{m_{\tilde{w}}^2}{m_{s_t}^2} \left(\frac{P \cdot \ell}{\tilde{w} \cdot \ell} + \frac{P \cdot \tilde{w}}{\tilde{w} \cdot \ell} - m_{\tilde{w}}^2 \frac{P \cdot \ell}{(\tilde{w} \cdot \ell)^2} \right) - 1 \\ b &= \frac{m_{\tilde{w}}^2}{m_{s_t}^2} \left(\frac{3}{2} \frac{P \cdot \ell}{\tilde{w} \cdot \ell} + \frac{P \cdot \tilde{w}}{\tilde{w} \cdot \ell} - m_{\tilde{w}}^2 \frac{P \cdot \ell}{(\tilde{w} \cdot \ell)^2} \right) - \frac{P \cdot \tilde{w}}{m_{s_t}^2} - \frac{1}{2} + \lambda a \\ c &= \frac{m_W^2 - 2\tilde{w} \cdot \ell}{m_{s_t}^2} \left(2 \frac{P \cdot \ell}{m_{\tilde{w}}^2} - 2 \frac{\tilde{w} \cdot \ell}{m_{\tilde{w}}^2} + \frac{P \cdot \tilde{w}}{\tilde{w} \cdot \ell} - m_{\tilde{w}}^2 \frac{P \cdot \ell}{(\tilde{w} \cdot \ell)^2} \right) + \\ &\quad + 2\lambda b - \lambda^2 a \end{aligned} \quad (15)$$

and

$$\lambda = (2\tilde{w} \cdot \ell - m_{\tilde{w}}^2 - m_W^2) / m_{\tilde{w}}^2 \quad (16)$$

Here $x = \sqrt{x_T^2 + x_L^2} = 2p \cdot \ell / (m_{s_t}^2)$, $y = 2\tilde{w} \cdot \ell / (m_{\tilde{w}}^2)$, and p , $\tilde{w}$, and ℓ are the four-momenta of s_t , $\tilde{w}$, and ℓ^+ , respectively. Similarly the charged lepton p_T distribution arising from t quark decay is given by (exactly in $\varepsilon = m_t^2/m_W^2$):

$$\frac{dN}{dx_T} = \frac{12}{f(\varepsilon)} \int_{x_T}^1 \frac{dx}{\sqrt{x^2 - x_T^2}} \frac{x(1-x)}{1-\varepsilon x} \quad (17)$$

where

$$f(\varepsilon) = \frac{1}{\varepsilon^3} \left(12 - 6\varepsilon - 2\varepsilon^2 + 12 \frac{1-\varepsilon}{\varepsilon} \ln(1-\varepsilon) \right) \quad (18)$$

In fig. 4 the p_T spectrum of the charged lepton from s_t decay is compared with the corresponding spectrum from t decay, for different $\tilde{w}$ masses. It can be seen that the two spectra are not too different even for a very light $\tilde{w}$, in which case the lepton spectrum from s_t is softer than that from t . For heavier $\tilde{w}$ they become essentially identical. We also studied the transverse mass distributions, with $M_T^2 = [(\sum_i |\vec{p}_{T_i}|)^2 - (\sum_i \vec{p}_{T_i})^2]$. The apparent scalar top transverse mass peaks at a substantially smaller value for the mass than the true value (fig. 5). This is due to the presence of two neutrino-like particles in the final state, so that a bias is introduced by treating the system of two particles as a single massless object, as one would do in a top quark search.

For $m_{s_t} \simeq 40$ GeV and $40 \text{ GeV} \lesssim m_{3/2} \lesssim 150$ GeV all diagrams in fig. 3 may be important. Final states with an additional hadron jet ($\tilde{g}$) associated with the charged lepton may arise. Otherwise the main qualitative conclusions of the previous discussions are unchanged. For a moderate $\tilde{g}$ mass, $B_{SL}(s_t)$ is still competitive with $B_{SL}(t)$.

In conclusion, if candidate events for top quark-antiquark associated production at the Collider are identified, it has still to be decided whether the top quark or the top scalar has been discovered. In fact we have argued that the production rates for a scalar top and for a top quark of equal masses are not very different and are almost the same within the uncertainty on both. The semileptonic branching ratios can also be rather close. Furthermore, although the lepton p_T and the transverse mass distributions are somewhat different, they do not provide a simple way of distinguishing between the two possibilities, as they would only reflect in a different mass determination. The most direct way of telling a scalar top from a top quark is to look at the W decay products. It is in fact quite unlikely that a scalar top will be found in W decays, as opposed to the top quark.

REFERENCES

- [1] For a recent review see, for example, S. Ferrara, *in* Proc. Int. Europhysics Conf. on High-Energy Phys., Brighton, 1983 (Rutherford Appleton Lab., 1983), p. 522.
- [2] E. Cremmer, S. Ferrara, L. Girardello and A. van Proyen, Phys. Lett. **116B** (1982) 231.
R. Barbieri, S. Ferrara and A. Savoy, Phys. Lett. **119B** (1982) 343.
R. Arnowitt, A.H. Chamseddine and P. Nath, Phys. Rev. Lett. **49** (1982) 970 and **50** (1983) 232.
S. Ferrara, D.V. Nanopoulos and C. Savoy, Phys. Lett. **123B** (1983) 214.
E. Cremmer, P. Fayet and L. Girardello, Phys. Lett. **122B** (1983) 41.
H.P. Nilles, M. Srednicki and D. Wyler, CERN-TH 3432 (1982).
N. Ohta, Tokyo UT-3881 (1982).
- [3] L.E. Ibanez, Phys. Lett. **118B** (1982) 73.
L.E. Ibanez and C. Lopez, Phys. Lett. **126B** (1983) 54.
J. Ellis, D.V. Nanopoulos and K. Tamvakis, Phys. Lett. **121B** (1983) 123.
L. Alvarez-Gaumé, J. Polchinski and M.B. Wise, Nucl. Phys. **B221** (1983) 495.
J. Ellis, J. Hagelin, D.V. Nanopoulos and K. Tamvakis, SLAC-PUB-3042 (1983).
J. Ellis and S. Rudaz, Phys. Lett. **128B** (1983) 248.
- [4] P. Fayet, Phys. Lett. **125B** (1983) 178.
- [5] A. Dicus, J. Nandi and X. Tata, Phys. Lett. **129B** (1983) 451.
P. Nath, R. Arnowitt and A.H. Chamseddine, Harvard Univ. preprint HUTP-83/A077 (1983).
- [6] S. Weinberg, Phys. Rev. Lett. **50** (1983) 387.
- [7] M. Veltman, Nucl. Phys. **B123** (1977) 89.
- [8] R. Barbieri and L. Maiani, Nucl. Phys. **B224** (1983) 32.
- [9] For a recent review see, for example, F. Halzen, Proc. Int. Conf. on High-Energy Physics, Paris, 1982 [J. Phys. **43** (1982) C3-381].
- [10] P.R. Harrison and C.H. Llewellyn-Smith, Nucl. Phys. **B213** (1983) 223, **E B223** (1983) 542.
See also I. Antoniadis, L. Baulieu and F. Delduc, PAR-LPTHE 83-10 (1983).
- [11] B.L. Combridge, Nucl. Phys. **B151** (1979) 429.
V. Barger, F. Halzen and W.Y. Keung, Phys. Rev. **D24** (1981) 1428 and **D25** (1982) 112.
R. Odorico, Phys. Lett. **107B** (1981) 231.
P. Mazzanti and S. Wada, Bologna preprint IFUB 82/9 (1982).
- [12] J. Ellis, J.S. Hagelin, D.V. Nanopoulos and M. Srednicki, Phys. Lett. **127B** (1983) 233.
B. Grinstein, J. Polchinski and M.B. Wise, Phys. Lett. **130B** (1983) 285.
- [13] G. Altarelli, B. Mele and S. Petrarca, CERN-TH 3822 (1984).

Figure captions

- Fig. 1 Mass spectrum of SUSY particles, in terms of the gravitino mass $m_{3/2}$, in the schematic framework explained in the text, for $m_{s_t} \simeq 40$ GeV, where m_{s_t} is the mass of the scalar top. The experimentally allowed range of masses lies between the shaded boundaries.
- Fig. 2 Comparison of the transverse momentum distributions of the top quark and the scalar top from gg fusion and $q\bar{q}$ annihilation in $p\bar{p}$ collisions.
- Fig. 3 Diagrams contributing to scalar top decays: f_U and f_D stand for general up and down quarks and leptons in a weak doublet, and s_U, s_D are the associated SUSY scalars.
- Fig. 4 Transverse momentum distributions of the charged lepton from $s_t \rightarrow b\tilde{w}^+ \rightarrow b\ell^+ \nu \tilde{\gamma}$ (full curves) and $t \rightarrow b\ell^+ \nu$ (dashed curve).
- Fig. 5 The transverse mass distributions for s_t and t semileptonic decays. For s_t the final ν and $\tilde{\gamma}$ have been mistreated as a single massless object.

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5