Supersymmetry and Charged Current Events at HERA

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Abstract

A light stop, with an R-parity-violating coupling λ'_{131}, has been suggested as an explanation of the excess in high-Q^2 neutral current events observed at the HERA collider. We show that in this scheme a corresponding excess in charged current events — such as that reported by the H1 Collaboration — can appear naturally, without calling for the presence of light sleptons or additional R-parity-violating couplings, if there exists a chargino lighter than the stop. The predicted event shapes agree well with the data. The relevant region of parameter space is identified, taking into account constraints coming from precision electroweak measurements, atomic parity violation and recent searches for first-generation leptoquarks at the Tevatron collider.

The reported [1,2] excess of neutral-current (NC) events at large Q^2 in deep inelastic scattering (DIS) at the HERA collider has been interpreted as a possible hint of physics beyond the Standard Model (SM) [3–12]. Among the various solutions considered, the currently favoured ones involve the s-channel production of a scalar 'leptoquark', which couples to the positron and one of the valence (or sea) quarks in the proton. Scalars of this nature, namely *squarks*, form an essential ingredient of R-parity violating ($\not{R}_p$) supersymmetric theories. Supersymmetry (SUSY) not only provides one of the most well-motivated extensions of the SM but also leads to an explanation of the observed NC excess [4–7,12]. Its role as a framework for solutions to the HERA anomaly thus merits serious consideration.

Within the framework of the minimal supersymmetric extension of the Standard Model (MSSM) the relevant $\not{R}_p$ contribution to the superpotential reads

$$\mathcal{W} = \lambda'_{ijk} L_i Q_j D_k^c \; , \tag{1}$$

where i, j and k are generation indices, L_i and Q_j are the $SU(2)_L$ doublet lepton and quark superfields respectively, while D_k^c is a charge-conjugate right-handed down quark superfield. In view of the various constraints on $\not{R}_p$ couplings from low-energy processes [13], the most natural interpretation of the HERA NC data is given by the s-channel production [14] of a left-handed charm ($\tilde{c}_L$) or top ($\tilde{t}_L$) squark, with a mass of $\mathcal{O}(200)$ GeV and a coupling to the *valence d*-quark,

$$\lambda'_{1j1} \sim \frac{0.04}{\sqrt{\beta_{ed}}}, \qquad \beta_{ed} \equiv \mathrm{Br}(\tilde{q}_j \to e^+ d) \; , \tag{2}$$

where $j = 2, 3$. It is perhaps worth mentioning that a sizeable effect proceeding from the interaction of the positron with the *sea* quarks in the proton would, in general, demand rather large couplings, which are generally inconsistent with constraints from low-energy data. The only exception is the possibility [5,6] of stop production from a sea s-quark through $\lambda'_{132} \simeq 0.4/\sqrt{\beta_{ed}}$. The parameter space in this scenario has subsequently been shown to be constrained by LEP measurements [15] as well as by limits on the electron-neutrino mass [16]. In the following, we shall concentrate on a valence quark interpretation of the HERA anomaly.

Searches at the Tevatron collider for generic leptoquarks decaying into lepton and jet—of which squarks of $\not{R}_p$ supersymmetry are one example—put strong restrictions on the allowed branching ratio β_{ed}. For a squark mass of $\mathcal{O}(200)$ GeV, the bounds from the CDF [17] and D0 [18] analyses can be combined to indicate

$$\beta_{ed} \lesssim 0.5 \; . \tag{3}$$

While this result seems to rule out most explanations of the HERA anomaly involving leptoquark resonances[1], which have no other decay channels, supersymmetric models can easily get around the impasse. Given the rich particle spectrum of SUSY models, R-parity-*conserving* interactions can induce additional decay processes (not counting, of

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[1]Attempts have been made, however, to construct models that evade this constraint [19].

course, those that lead to hard charged leptons plus jets in the final state). If this is indeed the case, with the negative results at the Tevatron being due to an R-parity-conserving decay that is unobservable with the current data sample [20], then it is natural to ask what is predicted at HERA, where a major fraction $(1 - \beta_{ed})$ of the produced squarks will decay through these modes. One should typically expect some additional signals beyond those expected within the SM.

As a matter of fact, the H1 Collaboration [1]—and recently, the ZEUS Collaboration [21] as well—*have* reported a small excess in charged current (CC) events at large Q^2 in the same data sample that contains the NC excess. Although less significant statistically than the latter, the observed number is nevertheless too large to be consistent with SM background calculations at 2.5 standard deviations. In view of the expectation that R-parity-conserving decays of the produced squark would be observed at HERA, an identification of the CC excess with these modes seems to be the most natural conclusion. It is thus interesting to ask whether SUSY scenarios can admit this possibility without contradicting other known phenomenological results. Such an investigation assumes particular importance since it is difficult [12] to accommodate the CC excess within other non-supersymmetric explanations for the NC events. Some attempts in this direction have already appeared in the literature [4, 12, 22, 23]. In Ref. [4], it was pointed out that R-parity-conserving decays of the squark could possibly lead to the CC signal. However, the observed event shape — low jet multiplicity, large missing momentum, relatively large hadronic invariant mass of the final state, and, of course, a high Q^2 — need to be taken into account in any serious pursuit of such an explanation. A more detailed study of the event shape has, in fact, been undertaken, mainly in the context of charm-squark production, in Ref. [12], where an explanation of the CC excess is obtained by postulating a light slepton-sneutrino pair, each of which decays into jets through $\not{R}_p$ couplings. In order to obtain a signal consistent with the data, especially with the low jet multiplicity[2], the sneutrinos must be rather light, $m_{\tilde{\nu}} \lesssim 80$ GeV, while the wino mass M_2 should not be much larger than 200 GeV. Similar requirements lead to a CC signal consistent with the experimental data in the context of top-squark production [12, 22]. Although this is a plausible explanation of the HERA anomaly, light sneutrinos and low values of M_2 lead to a sizeable sneutrino pair-production cross section at LEP2 [25] and consequently to observable rates for four-jet final states [26]. This last prediction presents a knotty problem, in view of the difference between results presented by the ALEPH Collaboration and the other LEP experiments [27], but the general hope is that a definitive result may be obtained from the next run of LEP2. It is also worth pointing out that the solution suggested in Ref. [12] may lead, for scharm production, to events with a small fraction of wrong-sign $\ell^- +$ jet final states at HERA, which could be detectable [28] as more data accumulate. A final state with $\mu^+ +$ jets (with or without missing momentum) can also be accommodated in their scenario in the presence of a light smuon-sneutrino pair.

In this letter, we describe a scenario in which a CC excess with the appropriate event shape may be generated even in the absence of light sleptons. We concentrate on the stop

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[2]Although, even then, the CC events would, in general, tend to have two visible jets (see Fig.2 of Ref. [12]).

interpretation with an $\not{R}_p$ coupling λ'_{131} consistent with Eq. (2). A light stop in the 200 GeV mass range is natural in models with radiative breaking of electroweak symmetry where squarks of the first two generations are heavy [29]. It is perhaps better motivated, therefore, than the competing option of a light charm-squark. The hierarchy of masses of the left-handed stop and the sbottom arising from the $SU(2)$-breaking condition,

$$m_{\tilde{t}_L}^2 - m_{\tilde{b}_L}^2 = m_t^2 - m_b^2 + m_W^2 \cos 2\beta > 0 \ , \tag{4}$$

suggests a simple way of generating CC events: if an R-parity-conserving decay mode of the stop leads to sbottoms in the final state, the CC events may be generated by the subsequent decay $\tilde{b}_L \to \bar{\nu}d$ of the sbottom through the same $\not{R}_p$ coupling λ'_{131}. (A similar scenario for the $\tilde{c}_L$ is not possible, since the $\tilde{s}_L$ would be somewhat heavier than the $\tilde{c}_L$ for $\tan\beta > 1$.)

In principle, if the sbottom is light enough ($m_{\tilde{b}} \lesssim 120$ GeV), a substantial branching fraction for the weak decay of a stop into a bottom squark and a W^+ (with the last going mainly into jets) may be allowed. This has been proposed recently by Kon *et al.* [23] as the cause of the CC excess at HERA. Although $\tilde{t} \to \tilde{b} + W^+ \to \not{p}_T +$ jets is a straightforward and economical scenario, it has several potential problems. First, it would mostly lead to CC events with a jet multiplicity of *two* or *three*, whereas, of the four events observed by H1, only one is consistent with more than a single jet [30]. Moreover, the hadronic mass distribution is rather broad and one has to appeal to the low statistics [23] to claim any agreement with the data. Apart from such kinematic arguments, one also encounters the indubitable fact that the contribution of the stop–sbottom doublet to the ρ parameter tends to be too large ($\Delta\rho_{\tilde{t}} \gtrsim 2 \times 10^{-3}$) if the stop–sbottom mass splitting is so large. Slightly smaller values of $\Delta\rho$ may be achieved by introducing large mixing angles in the stop and sbottom sectors, and tuning them appropriately. This has, in fact, been invoked by the authors of Ref. [23]. However, even then, $\Delta\rho_{\tilde{t}} \gtrsim 1.5 \times 10^{-3}$ and hence the scenario is only marginally consistent at the 2σ level [24]. In addition, constraints from atomic parity-violation experiments tend to be more severe in the case of large stop mixing [31].

All these difficulties can, of course, be overcome trivially by pushing the sbottom mass to larger values. This would immediately make the jets coming from the W^+ particle softer, and hence less often visible, while increasing the missing p_T and the invariant hadronic mass of the CC events. It also reduces the contribution to the ρ parameter. Unfortunately, when the sbottom mass is pushed high enough to reduce the contribution to the ρ parameter to acceptable values, the R-parity-conserving branching fraction becomes very much smaller than the $\not{R}_p$ one, and the corresponding CC event rate gets highly suppressed.

It is thus clear that a phenomenologically consistent explanation for the CC observation will require light supersymmetric particles in addition to the $(\tilde{t}, \tilde{b})$ pair. The light slepton option has been commented on above and will not be looked into any further. Since the stop decay into neutralinos is naturally suppressed[3], the obvious course is to have a

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[3]One could consider a non-negligible mixing between stop and the left-handed charm squark, and thus

chargino lighter than the stop. For most of the MSSM parameter space, this also means a neutralino that is considerably lighter and is, therefore, the lightest supersymmetric particle (LSP). Assuming that there are no other light sfermions, just three decay channels are now open to the chargino, namely into the neutralino ($\tilde{\chi}^+ \to \tilde{\chi}^0 W^* \to \tilde{\chi}^0 f\bar{f}'$), into the sbottom ($\tilde{\chi}^+ \to c\tilde{b}$) through the Cabibbo–Kobayashi–Maskawa (CKM) mixing (provided the sbottom is lighter than the chargino), and an $\not{R}_p$ decay through a virtual stop. For most of the parameter space of interest, however, the $\not{R}_p$ decay of the chargino has very small width. We shall assume, in the following, that the $\tilde{t}$–$\tilde{b}$ mass splitting is $\lesssim 35$ GeV, which seems optimum for the ρ-parameter constraint, and that $m_{\tilde{t}} > m_{\tilde{\chi}^+} > m_{\tilde{b}}$. Now, depending on the branching ratios, two distinct possibilities present themselves.

(*i*) *The chargino decays predominantly into the neutralino.* If the latter is lighter than the sbottom ($m_{\tilde{\chi}^0} < m_{\tilde{b}}$), it can have only a three-body decay ($\tilde{\chi}^0 \to b\bar{d}\nu_e,\ \bar{b}d\bar{\nu}_e$). In this case, the final state would typically have multijets and relatively low missing momentum. Such configurations will probably be lost in the DIS (CC) background coming from SM interactions. On the other hand, if the neutralino is heavier than the sbottom (which thus becomes the LSP), *i.e.*

$$m_{\tilde{t}} > m_{\tilde{\chi}^+} > m_{\tilde{\chi}^0} > m_{\tilde{b}}\ ,$$

then it would decay into $b + \tilde{b}_L$ and the $\not{R}_p$ two-body (and only) decay of this $\tilde{b}_L$ would result in a significant amount of missing momentum. Although, at first sight, it might seem that we would still have too many jets, most of these are now very soft as long as $m_{\tilde{t}} - m_{\tilde{b}} \lesssim 35$ GeV. A simple parton-level Monte Carlo simulation shows that most of the jets (apart from the one coming from the $\tilde{b}_L$ decay) fail to meet the minimum energy criteria of the H1 experiment for defining a jet [1,30]—this typically results in most events ending up with only a single 'jet' and almost none with more than two. Even when the parton energy is adequate, the second jet would be soft and broad, and it is debatable whether it would be detectable at all. It is quite easy to confirm that the kinematic profiles for these configurations match those for the observed CC events. Thus, a cascade decay of stop into chargino into neutralino into sbottom seems to be an attractive scenario. However, like the ones considered above, this too suffers from an inherent drawback. Since *all* the charginos produced in the stop decay (and we must remember that this is the dominant decay of the stop for $\beta_{ed} \lesssim 0.5$) are potential CC event candidates (modulo detection efficiency), we now tend to get too many of the latter. Unless one wishes to appeal to experimental effects and the low statistics, this number can be reduced only by providing the sbottom with another decay channel. What can that be? If we violate gaugino mass unification and choose the soft supersymmetry-breaking parameters such that the mass of the lightest neutralino ($\tilde{\chi}_1^0$) is well below that of the sbottom while the next-to-lightest neutralino ($\tilde{\chi}_2^0$) lies just below the chargino and is the one involved in the above chain, then, indeed, a new decay mode $\tilde{b}_L \to b\tilde{\chi}_1^0$ can be opened. A scan of the parameter space shows that this rather complicated scenario can be achieved for a limited part of the parameter space and that too, only if there are large deviations from gaugino

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allow the tree-level decay $\tilde{t}_L \to c\tilde{\chi}^0$. It can be shown, however, that the mixing angle must be very large (~ 1) to lead to the observed signal; such a large mixing angle is difficult to accommodate within any reasonable theoretical framework.

mass universality. Though this scheme works — as we have checked — it has limited aesthetic appeal and will not concern us any further.

(*ii*) *The chargino decays predominantly into the sbottom and a charm quark.* This mode is naturally suppressed by the smallness of the CKM matrix element $V_{cb} = 0.036 - 0.046$. Although some enhancement can be obtained in the event of a nonzero $\tilde{b}$–$\tilde{s}$ mixing, we prefer to be conservative and do not consider this option. Thus for $\tilde{\chi}^+ \to c\tilde{b}_L$ to be the dominant decay mode, the chargino–neutralino coupling which drives the competing decay ($\tilde{\chi}^+ \to \tilde{\chi}^0 f f'$) has to be quite small. Fortunately, this occurs naturally when $|\mu|$ is large compared to M_2. The lightest neutralino is then mostly bino-like and couples very weakly to the W, while the lighter chargino and the next-to-lightest neutralino are mostly wino-like and nearly degenerate in mass. Thus, the chargino decay into the second neutralino will be strongly suppressed, if not disallowed. The mass hierarchy obtained in this scheme is of the form

$$m_{\tilde{t}} > m_{\tilde{\chi}^+} \simeq m_{\tilde{\chi}^0_2} > m_{\tilde{b}} > m_{\tilde{\chi}^0_1} \ .$$

We find that a branching ratio $\beta_{ed} \lesssim 0.5$ can naturally be obtained for values of $M_2 \sim m_{\tilde{\chi}^+} \sim 190$ GeV) and corresponding values of λ'_{131} as determined by Eq. (2). In order that the jet resulting from the charm quark should not be observable, we also require the sbottom to be fairly heavy, typically in the 170–180 GeV range, which is already indicated by the ρ parameter constraint. The $\not{R}_p$ decay of the sbottom into $\bar{\nu} d$ now provides a hard jet and missing momentum, as in the cases discussed above, which can be the origin of the observed CC excess.

For this last scenario—which will concern us in the rest of this letter—we observe that the sbottom can also decay in the R-parity-conserving mode $\tilde{b}_L \to b\tilde{\chi}^0_1$, with the neutralino then decaying into

$$\tilde{\chi}^0_1 \to \bar{b} d \bar{\nu}_e \ , b \bar{d} \nu_e \ . \tag{5}$$

through an intermediate sbottom. If the neutralino is light enough, $m_{\tilde{\chi}_1} \sim 100$–140 GeV, the missing p_T of this decay mode will tend to be low, and will mostly fail the CC selection criteria imposed by the H1 Collaboration (see below). Potential CC events resulting from this decay channel will then be lost in the SM DIS background. We can thus have a reduced number of CC events without making any assumption beyond that of a large $|\mu|$ and $M_2 \sim 190$ GeV. The simplicity and economy of this scheme is apparent. In effect, we only assume:

A. a small stop–sbottom mass-splitting, which is easily achieved with a modicum of left–right mixing in the stop sector;

B. a chargino, which is mainly wino-like, lying between the stop and sbottom in mass; and

C. a light neutralino, which is dominantly bino-like, and whose mass is considerably below that of the sbottom, typically 100–140 GeV.

With these assumptions, a study of the predicted event shapes and efficiencies is now in order.

In the presence of left–right mixing ($\theta_{\tilde{t}} = 0$ implies that the lighter eigenstate is a pure $\tilde{t}_L$), the stop production cross section at HERA is given by

$$\sigma(m_{\tilde{t}}) = \frac{\pi}{16 E_e^0 E_p^0} \left(\lambda'_{131} \cos\theta_{\tilde{t}}\right)^2 \; K \; f_d \left(\frac{m_{\tilde{t}}^2}{4 E_e^0 E_p^0}\right) \; , \tag{6}$$

where E_e^0, E_p^0 are, respectively, the initial electron and proton energies in the laboratory frame and $K \simeq 1.3$ parametrizes the QCD correction [32]. For $f_d(x)$, the d-quark distribution within the proton, we use the MRS(H) [33] structure functions, calculated using the package PDFLIB [34]. In this analysis, we neglect effects due to initial state radiation. The NC and CC signals are estimated using a parton-level Monte Carlo event generator. For jets within the angular coverage of their detector, the H1 Collaboration requires a minimum transverse momentum of

$$p_T^{\text{jet}} > 15 \, \text{GeV} \; . \tag{7}$$

In our parton-level analysis, each final-state quark is assumed to give rise to jets. All such 'jets' that lie within a cone

$$\Delta R_{jj} \equiv \sqrt{\Delta\eta^2 + \Delta\phi^2} \leq 1.0 \tag{8}$$

of each other are 'merged' to form a single jet by the simple expedient of adding their four-momenta together vectorially and identifying the sum with the momentum of the combination. Here $\Delta\eta$ is the difference of pseudorapidities and $\Delta\phi$ denotes azimuthal separation.

For NC events, the H1 analysis requires that the final-state positron must lie within

$$10° < \theta_e < 145° \; , \tag{9}$$

(the proton direction is positive) and it must have a minimum transverse energy:

$$p_{Te} \geq 25 \, \text{GeV} \; . \tag{10}$$

In addition, the missing transverse momentum must satisfy $\not{p}_T / \sqrt{p_{Te}} \leq 3 \, \text{GeV}^{1/2}$. The DIS Lorentz invariants can be estimated using either of the pairs $(\theta_e, \theta_{\text{jet}})$ and (θ_e, E_e); the latter give :

$$y_e = 1 - \frac{E_e}{E_e^0} \sin^2\theta_e, \quad Q_e^2 = \frac{p_{Te}^2}{1 - y_e}, \quad M_e = \sqrt{\frac{Q_e^2}{y_e}},$$

where $E_e^0 = 27.5 \, \text{GeV}$ is the energy of the incoming positron beam. The NC events are required to satisfy

$$0.1 < y_e < 0.9 \; . \tag{11}$$

With the above set of cuts (and other standard ones required to remove cosmic-ray backgrounds) the H1 Collaboration [1] reports 12 events with $Q_e^2 > 15000\,\mathrm{GeV}^2$ for an integrated luminosity of $14.2\,\mathrm{pb}^{-1}$ against the expected DIS background of 4.7 ± 0.76 events.

For the CC events, on the other hand, the cuts (9)–(11) are no longer operative. Instead, the H1 analysis demands that the missing transverse momentum be large:

$$\not{p}_T = \sum p_T(\text{hadrons}) > 50\,\mathrm{GeV} \ . \tag{12}$$

Analogous to the NC case, one can define DIS variables [35] using only measurable hadronic variables:

$$y_h = \frac{\sum(E - P_{\|})}{2E_e^0}, \qquad Q_h^2 = \frac{\not{p}_T^2}{1 - y_h} \ ,$$

where the sum runs over all observed hadronic clusters (in our simulation, 'jets' that pass the cuts). With a selection criterion of

$$y_h < 0.9 \ , \tag{13}$$

the H1 Collaboration observes 4 events above $Q_h^2 > 15000\,\mathrm{GeV}^2$, whereas the expected SM background is 1.77 ± 0.87 events.

It is easy to see that the efficiency for the NC events should be solely determined by the stop mass $m_{\tilde{t}}$, and for $m_{\tilde{t}} \sim 200\,\mathrm{GeV}$, can be approximated by

$$\epsilon_{NC} \approx 0.56 \ , \tag{14}$$

where we assume detector efficiency to be unity. The CC efficiencies, on the other hand, *can* depend on all the relevant masses in the decay chain. Again, our computations show that the only significant dependence is on the bottom-squark mass, and

$$\epsilon_{CC} \approx 0.37\,(0.45) \pm 0.008, \qquad \text{for } m_{\tilde{b}} = 170\,(180)\,\mathrm{GeV} \ . \tag{15}$$

As the NC event topology has been studied elsewhere [4–7, 12], and shown to be consistent with the data, we shall not repeat the exercise here. In Fig. 1 we present the kinematic distributions for the CC events. To be concrete, we have chosen to depict it for $m_{\tilde{t}} = 200\,\mathrm{GeV}, m_{\tilde{\chi}^+} = 190\,\mathrm{GeV}$ and for three typical sbottom masses. It turns out that the phase-space distribution is not very sensitive to the exact value of $m_{\tilde{\chi}^+}$ and is primarily governed by the sbottom mass.

Each of Fig. 1(a–c) depict the p_T distribution of the three potential 'jets' for a given $m_{\tilde{b}}$. As expected, the hardest jet emanates mainly from the $\tilde{b}$-decay, but with (small) contributions from the soft jets that merged with it. Consequently, its p_T-distribution just reflects the Jacobian peak with a slight smearing due to the small transverse boosts. For the next-to-hardest jet (dashed lines), the situation is quite different. In general, it has relatively low p_T and is mostly rejected by the cut (7). However, depending on the $\tilde{t}$–$\tilde{\chi}^+$ and the $\tilde{\chi}^+$–$\tilde{b}$ mass splittings, some configurations may have sufficient p_T to be
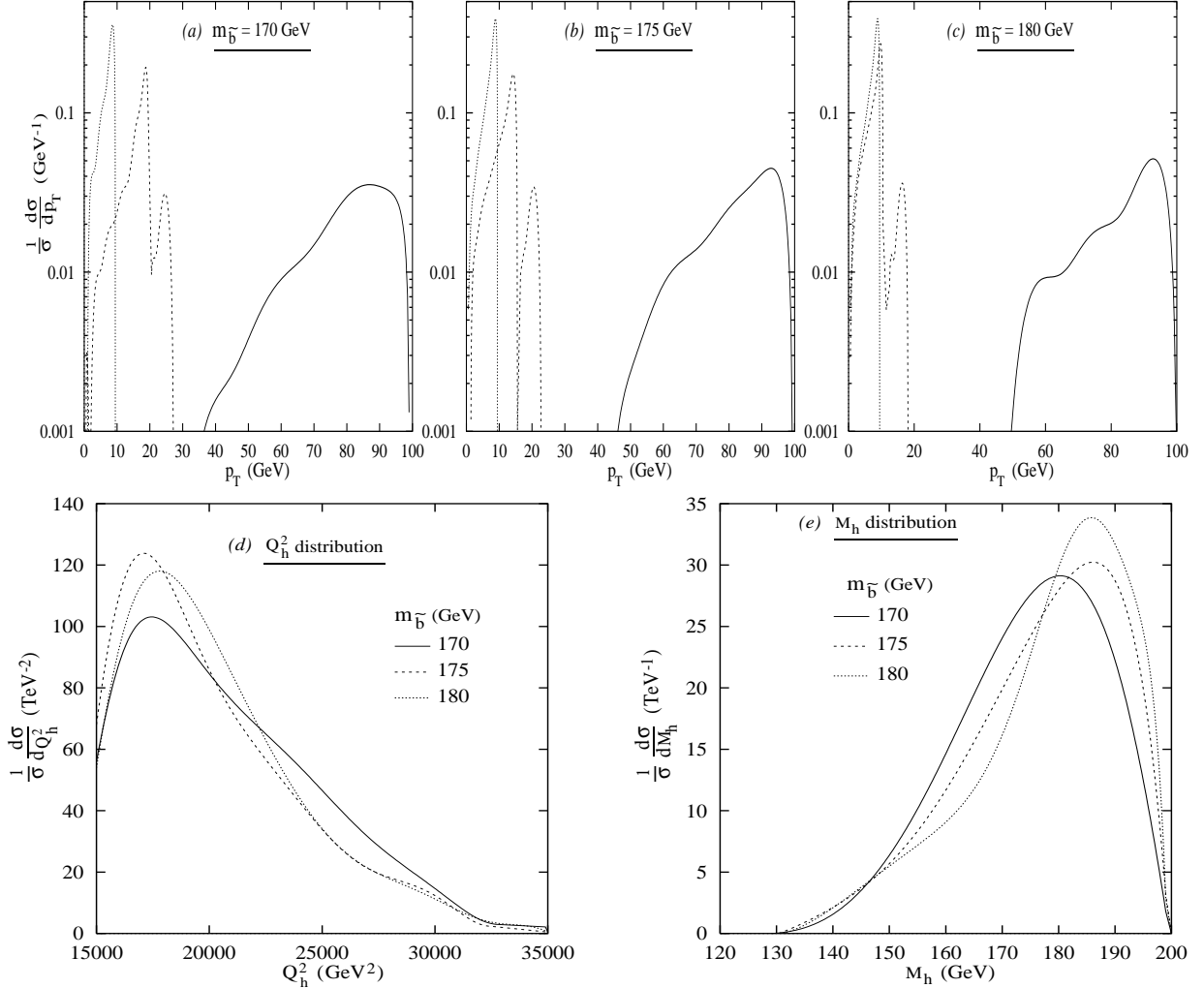
Figure 1: *Event topology for $m_{\tilde{t}} = 200$ GeV, $m_{\tilde{\chi}^+} = 190$ GeV, and for different values of $m_{\tilde{b}}$. (a–c) The p_T distributions for the three 'jets' in each case. (d) The Q_h^2 and (e) the M_h distributions for all three cases.*

observable. The remaining jet, if it has not merged with the hardest two, is always too soft to be seen. This is amply demonstrated by the dotted lines. Thus, even though we started with as many as three quarks in the final state, we never have more than two jets at the observable level. Moreover, the fraction of 2-jet events is always much smaller than that of single-jet events for $m_{\tilde{b}} \sim 180$ GeV; the fraction naturally goes up as the sbottom mass decreases. For the case where two jets can indeed be seen, their angular separation tends to be evenly distributed within the interval $1 \leq \Delta R_{jj} \leq 3$.

Figure 1(d) shows that the number of CC events is a rather sensitive function of Q_h^2; most of the events are concentrated within $15000 \text{GeV}^2 < Q_h^2 < 25000 \text{GeV}^2$. This is well in agreement with the present HERA data [1,21]. Events with $Q_h^2 > 30000 \text{GeV}^2$ [21] would, however, be harder to explain within this scenario. The mass distribution in Fig. 1(e) peaks at about 180–190 GeV, but has a long tail, which may account for the fact that one of the observed events has a low M_h (~ 157 GeV). It is interesting that such a peak in the mass distribution can be obtained in the scenario of Ref. [12] only for rather small

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sneutrino masses, while the scenario of Ref. [23] has a much flatter and broader mass distribution.

In Fig. 1 we present results for only a few points in the available parameter space, but it is easy to see that the same qualitative features would hold for other points as long as these satisfy the conditions (A-C) above. What of the events resulting from the decay chain of Eq.(5)? A simulation analogous to the one described above shows that most of these events indeed fail to survive the cuts of Eqs.(7, 12–13), with the missing momentum cut proving to be the most severe. The corresponding efficiency is a sensitive function of $m_{\tilde{\chi}^0}$ and ranges from about 4% to 13% as the latter varies from 90 GeV to 140 GeV. This is rather too low to yield an observable number of multijet events with large missing momentum in the present data sample.

Having established that the HERA results for CC events can be explained within the present scenario, it becomes interesting to identify the part of the MSSM parameter space that supports our solution. Certain requirements are immediately obvious. As explained above, since we want the chargino to be wino-like and slightly lighter than the stop, we are immediately constrained to $M_2 \sim 190$ GeV. For μ, $\tan\beta$ and the ratio M_1/M_2, the situation is more complicated. We note that the ratio of the number of NC events to CC events in this scenario is given by

$$\frac{N_{CC}}{N_{NC}} = \left(\frac{1 - \beta_{ed}}{\beta_{ed}}\right) \, \text{Br}(\tilde{\chi}^+ \to c\tilde{b}) \, \text{Br}(\tilde{b} \to \bar{\nu}d) \, \left(\frac{\epsilon_{CC}}{\epsilon_{NC}}\right) \, , \tag{16}$$

with efficiencies as in Eqs. (14) and (15). Consider a typical value $\beta_{ed} \simeq 0.4$ consistent with the Tevatron constraints. Since $\text{Br}(\tilde{b} \to \bar{\nu}d) \approx 0.4$ (0.6) for $m_{\tilde{\chi}^0} \sim 100$ (140) GeV, we have $N_{CC}/N_{NC} \simeq (0.35\text{–}0.5) \times \text{Br}(\tilde{\chi}^+ \to c\tilde{b})$. A good agreement with the H1 rates thus requires $\text{Br}(\tilde{\chi}^+ \to c\tilde{b}) \gtrsim 0.5$. This constrains the parameter space in the μ–$\tan\beta$ plane, and we present the favoured region in Fig. 2. While this region depends on the value of the $\not{R}_p$ coupling λ'_{131} as well as on the squark masses (these not only determine the production cross section and the efficiency, but also bound M_2), we present the parameter space for only one representative set of values. To be deemed acceptable, a point in the parameter space was required to lead to 7 ± 1 NC events and 2.5 ± 1 CC events for an integrated luminosity of 14.2 pb^{-1} and simultaneously satisfy $0.3 \leq \beta_{ed} \leq 0.5$. Note that these bounds on β_{ed} are relaxed for a slightly heavier stop (say 210 GeV). This would considerably expand the favoured region. A similar effect can be obtained by considering a mixing the $\tilde{b}$–$\tilde{s}$ sector. Our estimate should thus be regarded as a conservative one.

As shown by the solid lines in Fig. 2, for our scenario to work within the scheme of universal gaugino masses ($M_1/M_2 = \frac{5}{3}\tan^2\theta_W \simeq 0.5$), a relatively large $|\mu|$ is indicated. The curious shape of the curves owes its origin to the existence of minima in the wino content of the LSP as $\tan\beta$ is varied. The asymmetry between positive and negative μ can be traced to a similar source. While $\mu > 1$ TeV (or $\mu < -600$ GeV) is perfectly acceptable, it is interesting to ask if smaller values could be accommodated. An elegant way to achieve this is to reduce the mass difference $m_{\tilde{\chi}^+} - m_{\tilde{\chi}^0}$ so that the two-body decay $\tilde{\chi}^+ \to \tilde{\chi}^0 W^+$ is suppressed kinematically and the relevant process is the *three-body* decay

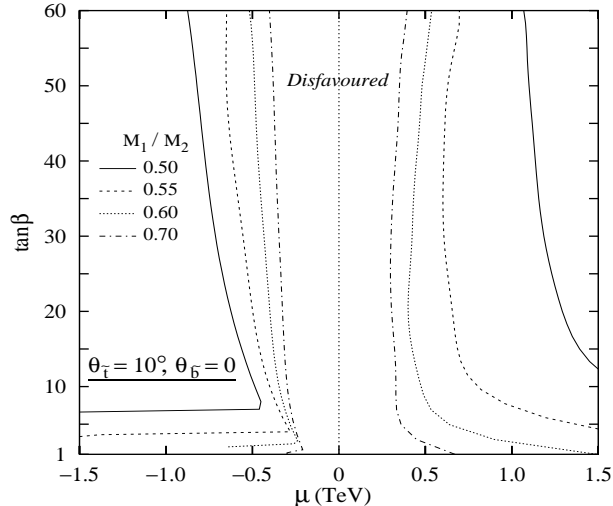Figure 2: *The favoured range (see text) in the μ–$\tan\beta$ plane for $(\lambda'_{131} = 0.07, m_{\tilde{t}} = 200\,\mathrm{GeV}, m_{\tilde{b}} = 175\,\mathrm{GeV})$ for four values of M_1/M_2. We assume a small stop mixing.*

of the chargino into the LSP and two light fermions. This happens naturally in the event of small corrections to the universality relation in the gaugino sector. In Fig. 2, we also show the favoured region for three other values of M_1/M_2. *A priori*, the dependence on this ratio could be quite complicated. While a larger value of M_1/M_2 would kinematically suppress the $\tilde{\chi}^+ \to \tilde{\chi}^0 f\bar{f}'$ channels, it would also increase the higgsino content of both the chargino and the LSP and consequently the $\tilde{\chi}^+\tilde{\chi}^0 W$ coupling. At the same time, the $\tilde{\chi}^+ c\tilde{b}$ coupling decreases too. Still, for the range of interest, the dependence is monotonic. It must be noted though that $M_1/M_2 > 0.7$ implies $m_{\tilde{\chi}^0} \gtrsim 140\,\mathrm{GeV}$, which runs the risk of resulting in a few (possibly observable) multijet events with large missing momentum.

In Fig. 2, we have chosen particular values for the squark mixing. A non-zero value for $\theta_{\tilde{b}}$ will serve to modify both $\mathrm{Br}(\tilde{\chi}^+ \to c\tilde{b})$ and $\mathrm{Br}(\tilde{b} \to \bar{\nu}d)$. Since the $\tilde{b}\bar{b}\widetilde{B}$ coupling vanishes exactly for $\tan\theta_{\tilde{b}} = 0.5$, for a mixing angle in the vicinity of this value, the curves in Fig. 2 would move *inwards* by approximately 150 GeV. Stop mixing, on the other hand, manifests itself mainly in determining the stop-chargino coupling and hence β_{ed}. For $|\theta_{\tilde{t}}| \lesssim 15°$, the dependence on this parameter is negligible for negative μ. For positive μ, the deviation from Fig. 2 is maximum for $M_1/M_2 \sim 0.5$ and, for moderate $\theta_{\tilde{t}}$, could, at best, lead to a shift by ~ 100–150 GeV, with the weakest constraints being obtained for $\theta_{\tilde{t}} \approx 0$.

A further interesting feature of this scenario is the possibility of the chargino decaying into a neutralino and a $\ell^+\nu$ pair instead of jets. This would typically lead to a hard lepton, jets and missing momentum in the final state. For $\ell = \mu$, this could perhaps explain [36] the muonic event [37] observed by H1. A similar event with the τ^+ leads to a narrow jet and is swamped by the CC DIS signal. An e^+ with low missing momentum would, similarly, be indistinguishable from the NC DIS, but events with larger missing momentum might be detectable as statistics multiplies.

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We now turn to the various low-energy measurements that, in principle, could further constrain the allowed parameter space. As we have seen in the above discussions, if the stop is responsible for both the NC and the CC excesses at HERA, it cannot be a pure left-handed state. The corrections to the ρ parameter would then be too large to be consistent with precision electroweak measurement data. While a large mixing can evade this problem, such a solution runs counter to the bounds from atomic parity violation [31]. The effective weak charge measured experimentally now receives a contribution from both stops; for a nucleus $^Z X_A$, it is given by

$$\Delta Q_W = -\frac{(\lambda'_{131}\cos\theta_{\tilde{t}})^2}{2\sqrt{2}G_F}(2A-Z)\left[\frac{1}{m_{\tilde{t}}^2}+\frac{\tan^2\theta_{\tilde{t}}}{m_{\tilde{T}}^2}\right]. \tag{17}$$

where $\tilde{T}$ denotes the heavier stop. The experimental bound [38] on ΔQ_W of the Cesium atom, along with the size of the NC excess at HERA, thus translates into a bound in the $(\theta_{\tilde{t}}, m_{\tilde{T}})$ plane. For example, $m_{\tilde{t}} \simeq 200$ GeV implies, at the 95% C.L. level,

$$\beta_{ed} \gtrsim 0.3\left[1+\tan^2\theta_{\tilde{t}}\,\frac{m_{\tilde{t}}^2}{m_{\tilde{T}}^2}\right].$$

It should be remembered though that the error in ΔQ_W is dominated by theoretical uncertainties, and so this bound should be considered with great caution. It is clear that, given the upper bound on β_{ed} imposed by the Tevatron data, small values of the mixing and/or a large hierarchy between the stop masses are preferred.

A simultaneous resolution of the ρ-parameter and the atomic parity-violation constraints thus call for a small $\tilde{t}$–$\tilde{b}$ mass splitting as well as a small stop mixing angle. While this may seem difficult at first, in reality, it can arise in a natural way. Consider the case where m_U, the supersymmetry-breaking mass parameter in the right-handed stop sector, is large in comparison with that in the left-handed sector (m_Q). The lightest stop mass is then approximately given by

$$m_{\tilde{t}}^2 \simeq m_Q^2 + D_L^t + m_t^2\left(1-\frac{\widetilde{A}_t^2}{m_U^2}\right), \tag{18}$$

where $\widetilde{A}_t = A_t - \mu/\tan\beta$ and $D_L^t = 0.35 M_Z^2 \cos 2\beta$. The lighter stop will be mainly left handed (*i.e.* $\theta_{\tilde{t}}$ is small), and the heavier stop mass ($m_{\tilde{T}} \sim m_U$) may be pushed to large values. Clearly, the smaller $\theta_{\tilde{t}}$ is, the higher $m_{\tilde{T}}$ is pushed—see Fig. 3(a)—and consequently, the smaller the additional contribution to ΔQ_W is. Ignoring the small mixing in the sbottom sector, the left-handed sbottom mass will be given by

$$m_{\tilde{b}}^2 \simeq m_Q^2 + D_L^b + m_b^2\ , \tag{19}$$

where $D_L^b = -0.42 M_Z^2 \cos 2\beta$. It is clear that, under these conditions, for values of $|\widetilde{A}_t| \sim m_U$, the lightest stop and sbottom masses will be close and hence the contribution to the ρ parameter will be small. On the other hand, for a fixed positive value of $m_{\tilde{t}} - m_{\tilde{b}}$, the contribution to the ρ parameter is smaller for larger values of m_U and for $|\widetilde{A}_t| \sim m_U$. This conclusion remains valid as long as the mixing in the sbottom sector is small. Large
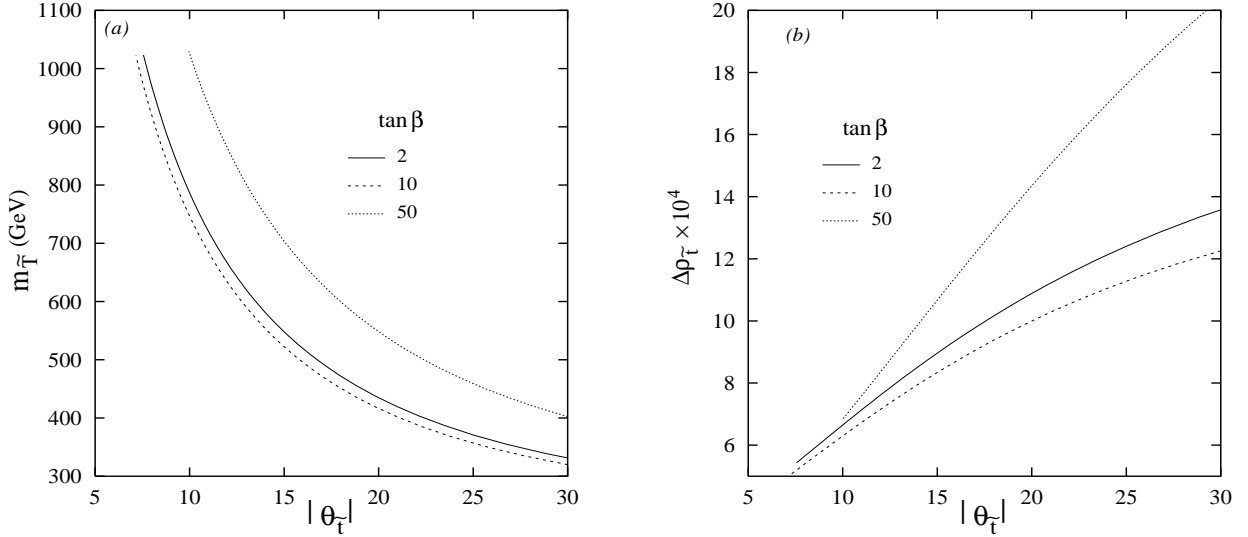
Figure 3: (a) *The heavier stop mass as a function of the mixing angle for* $m_{\tilde{t}} = 200$ *GeV,* $m_{\tilde{b}} = 175$ *GeV,* $|\mu| = 1$ *TeV,* $A_b = 0$, $m_D = 2$ *TeV and* $m_U \leq 1$ *TeV.* (b) *The corresponding contribution of the stop–sbottom doublet to the parameter* $\Delta\rho$.

$\theta_{\tilde{b}}$ tends to push the value of m_Q up, leading to a substantial mixing in the stop sector for the same values of the stop masses. Sbottom mixing appears naturally in the large $\tan\beta$ region and its effects are clearly seen in Fig. 3(a).

In Fig. 3(b), we plot the stop contribution [39] to $\Delta\rho$ as a function of $\theta_{\tilde{t}}$ for $m_{\tilde{t}} = 200$ GeV and $m_{\tilde{b}} = 175$ GeV. Fixing μ and m_D, the right-handed sbottom mass parameter, allows us to see the effects of a small mixing in the sbottom sector for large values of $\tan\beta$. Large values of μ and $\tan\beta$ also have a striking effect on the branching ratio $\mathrm{Br}(b \to s\gamma)$ as we shall discuss below. As argued above, the larger the left-handed component of the lighter stop, the smaller $\Delta\rho$. The experimental numbers suggest [24] $\Delta\rho_{\tilde{t}} < 1.5 \times 10^{-3}$ for $m_t = 175\,\mathrm{GeV}$ and $m_h \simeq 100\,\mathrm{GeV}$. From the figure, it is apparent that $\theta_{\tilde{t}} \lesssim 15°$ is preferred by the data. Observe that even lower values of $\Delta\rho$ may be obtained by relaxing the upper bound on m_U.

Since both stop masses are determined in this scenario, it is necessary to discuss their impact on the mass of the lightest CP-even Higgs [40]. In Fig.4(a), we present the variation of the Higgs mass as a function of the mixing angle, for the same values of the stop and sbottom masses. Although there are four branches associated with the two different signs of μ and the mixing parameter $\widetilde{A}_t$, for $m_A \gtrsim 250\,\mathrm{GeV}$, there is little difference between the branches. Hence, we plot the Higgs mass only for positive values of μ and $\widetilde{A}_t$. We see that small mixing angles are preferred for low values of $\tan\beta$, while no information on the mixing angle may be obtained from Higgs mass considerations in the large $\tan\beta$ regime.

Since the stop and chargino masses are fixed, we can determine the stop–chargino contribution to $\mathrm{Br}(b \to s\gamma)$ [41]. It is important to remember though that the physical
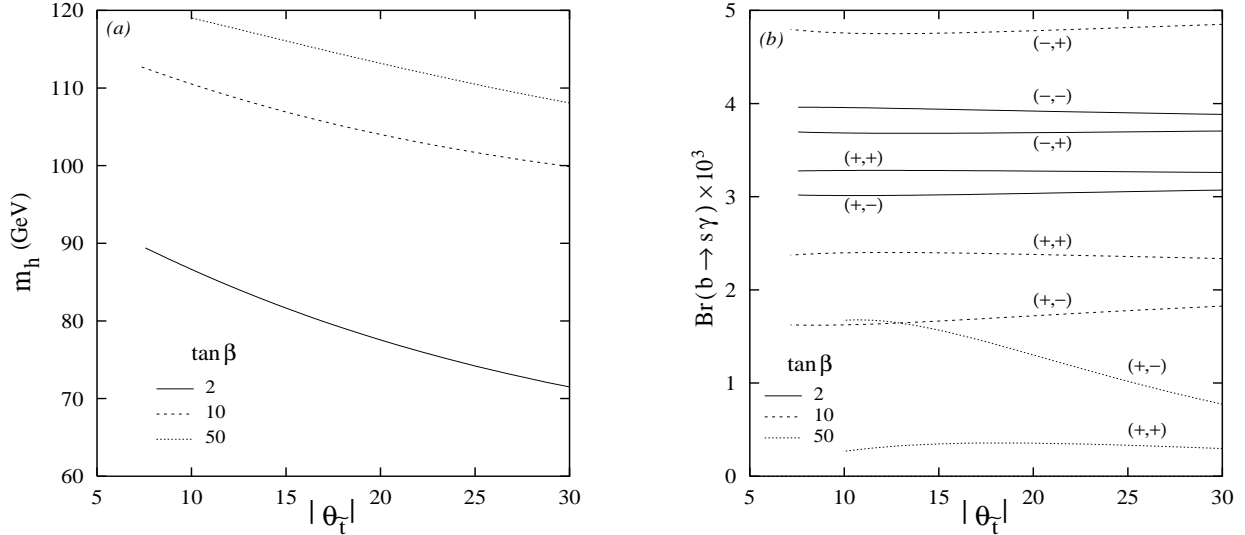
Figure 4: (a) *The mass of the lightest CP-even Higgs as a function of the stop mixing angle, for the same parameters as in Fig.3 and $m_A = 1$ TeV. (b) Br(b $\to$ sγ) as a function of the stop mixing angle for the same parameters. The branches are defined by* $(\mathrm{sgn}(\mu), \mathrm{sgn}(\widetilde{A}_t))$.

branching ratio depends on additional parameters which are not relevant for an interpretation of the HERA data. For a specific choice of these parameters and $m_{\tilde\chi^+} = 190$ GeV, Fig. 4(b), shows the branching ratio as a function of the stop mixing angle. We set the first and second generation squark masses to 2 TeV and consider all four branches defined by the signs of μ and $\widetilde{A}_t$. The QCD scale has been fixed at $Q = m_b/2$, a value that reproduces the next-to leading-order corrections[4] in the SM. Setting $Q = m_b$ would typically reduce the rates by about 20%, giving a measure of the theoretical uncertainties in these calculations. For $\tan\beta \gtrsim 10$, an acceptable rate may be obtained only by large cancellations between different contributions to the amplitude, which demands specific values of the Higgs and supersymmetric mass parameters. For the above parameter choice and $\mu = -1\,\mathrm{TeV}$, we find no solutions for $\tan\beta = 10, 50$. Similarly, for $\mu = +1\,\mathrm{TeV}$ and $\tan\beta = 50$, acceptable solutions may be found for only one of the branches of A_t. Although this seems to prefer lower values of $\tan\beta$, the situation can change for other parameter choices. In fact, with suitable parameters, we can always reproduce the experimental results for Br($b \to s\gamma$) in the scenario under consideration.

In all of the above, we have held $m_D = 2\,\mathrm{TeV}$. Although our scenario is consistent with lower values of this mass parameter, the analysis becomes more involved in the large $\tan\beta$ region, where the mixing in the sbottom sector becomes large. As noted earlier, a large value of $\theta_{\tilde b}$ will modify the relevant branching fractions, allowing for lower values of μ for a given M_1/M_2 and $\tan\beta$. On the other hand, since large $\theta_{\tilde b}$ generally implies a large $\theta_{\tilde t}$, the atomic parity violation constraints become more stringent in this case. Observe that, for an appropriate choice of the mass parameters, larger values of the sbottom mixing

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[4]A complete next-to leading-order computation in the supersymmetric model has not yet been performed.

angle may allow for a reduction of the value of $\Delta\rho$, even for moderately large values of the stop mixing angle.

To summarize, then, we have investigated the possibility that both the neutral and the charged current anomalies seen at HERA at high Q^2 are consequences of the resonant production of a stop, which decays to an e^+d final state to give the NC events and to $b\tilde{\chi}^+$ followed by cascade decays to give the CC events. Something of this nature is indicated by the negative results of the Tevatron search for leptoquarks/squarks in the 200 GeV mass range. Among the various scenarios discussed, the CC events are best explained through the decay $\tilde{\chi}^+ \to c\tilde{b}_L$, which is CKM-suppressed, but can be sizeable in the limit of a somewhat large $|\mu|$ parameter, when the competing decay channel is considerably reduced. The kinematic distributions in this scenario explain the observed CC excess rather well and the parameter space which supports this solution is perfectly consistent with electroweak precision measurements at LEP and with other constraints such as those originating from radiative B-decays, atomic parity violation and Higgs searches at LEP. Although it may seem premature to demand precise agreement with data for the CC excess reported by H1 and ZEUS, our scenario makes firm predictions which should be testable in forthcoming runs of the existing high-energy accelerator facilities.

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