Supersymmetry at the Electroweak Scale. *

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Abstract

The simplest interpretation of the global success of the Standard Model is that new physics decouples well above the electroweak scale. Supersymmetric extension of the Standard Model offers the possibility of light chargino and the right-handed stop (with masses below $M_Z$), and still maintaining the successful predictions of the Standard Model. The value of $R_b$ can then be enhanced up to $\sim 0.218$ (the Standard Model value is $\sim 0.216$). Light chargino and stop give important contribution to rare processes such as $b \to s\gamma$, $K^0 - K^0$ and $B^0 - \bar{B}^0$ mixing but consistency with experimental results is maintained in a large region of the parameter space. The exotic four-jet events reported by ALEPH (if confirmed) may constitute a signal for supersymmetry with such a light spectrum and with explicitly broken $R-$parity. Their interpretation as pair production of charginos with $m_C \sim 60$ GeV, with subsequent decay $C \to \tilde{t}_Rb \to dsb$ (where $m_{\tilde{t}} \sim 55$ GeV) leads to signatures very close to the experimental observations.

1. Introduction.

The Standard Model (SM) is a highly successful theory. The simplest interpretation of this fact is that any new physics decouples well above the electroweak scale. On the other side, there are the well known theoretical reasons (hierachy problem, ”naturalness” of elementary scalars etc.) to expect

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new physics to be close to the electroweak scale. Is this in contradiction with the success of the SM? Sypersymmetry is a particularly attractive extension of the SM. We are facing several interesting questions such as: how low the scale of supersymmetric particle masses can be to maintain consistency with the success of the SM? is there a room for some superpartners with masses at the electroweak scale? or, maybe there are even some hints for such a light spectrum? Recent important progress in the low energy supersymmetric phenomenology allows to address those questions in fully quantitative way, with conclusions which are very stimulating for further experimental search for supersymmetry at the LEP200 and the Tevatron!

The success of the SM is best measured by its description of the bulk of the electroweak data, with an excellent overall agreement. The only clear discrepancy is the present experimental value of $R_b = 0.2211 \pm 0.0016$ which is more than $3\sigma$ away from the theoretical prediction. The experimental value of $R_c$ is $1.6\sigma$ away from the prediction and this is statistically much less significant. Finally, there are two $\sim 2\sigma$ deviations in the leptonic left-right asymmetry and the parameter $A_b$. Both measurements come from SLAC and those deviations look merely like experimental problems of some mismatch between the SLAC and LEP data. Indeed, $A_{LR}^c$ is a measure of the $\sin^2 \theta^e_{\text{eff}}$ and it disagrees with the LEP measurements of this angle which can be extracted from the parameters $A_c, A_t, A_{FB}^t$. The direct SLAC measurement of $A_b$ disagrees with the indirect LEP determination from $A_{FB}^b$.

One remark is in order here. The SLD value of $A_{LR}^c$ gives $\sin^2 \theta^e_{\text{eff}} = 0.23049 \pm 0.00050$ whereas the LEP value is $\sin^2 \theta^e_{\text{eff}} = 0.23178 \pm 0.00031$ [1]. In the SM, the value of $\sin^2 \theta^e_{\text{eff}}$ can be very precisely calculated (instead of being determined from a global fit) in terms of $M_Z, m_t$ and $M_h$. We get e.g. the results shown in Table 1.

Table 1: Predictions in the SM for $\sin^2 \theta^e_{\text{eff}}$ for various top quark and Higgs boson masses. The error of this predictions (coming mainly from the uncertainty of the hadronic contribution to the photon vacuum polarization) is $\pm 0.00025$.

<table>
<thead>
<tr>
<th>$m_t$ (GeV)</th>
<th>170</th>
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<th>190</th>
<th>170</th>
<th>180</th>
<th>190</th>
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<tr>
<td>$M_h$ (GeV)</td>
<td>60</td>
<td>60</td>
<td>60</td>
<td>150</td>
<td>150</td>
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</tr>
<tr>
<td>$\sin^2 \theta^e_{\text{eff}}$</td>
<td>0.23135</td>
<td>0.23101</td>
<td>0.23066</td>
<td>0.23182</td>
<td>0.23149</td>
<td>0.23114</td>
</tr>
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</table>

It is clear that the SLD measurement favours larger (smaller) values of $m_t$ ($M_h$) ($m_t \approx (180 - 190)$ GeV, $M_h \approx 60$ GeV) which are larger (smaller) than those favoured by the LEP result for $\sin^2 \theta^e_{\text{eff}}$ ($m_t \approx 170$ GeV, $M_h \approx 150$ GeV). Such values give, however, worse fit to other electroweak observables ($M_W, \Gamma_Z, \sigma_h, A_{FB}^{b,h}$).
The precision of the data is already high enough to be sensitive to the Higgs boson mass (which enters into the calculations only logarithmically). The full SM fit (the fitted parameters are \(m_t\), \(M_h\) and \(\alpha_s(M_Z)\)) gives \(M_h = 76^{+93}_{-44}(1\sigma)^{+27}_{-76}(2\sigma)\). whereas in the fit without \(R_b\) and \(R_c\) we get \(M_h = 94^{+117}_{-55}(1\sigma)^{+346}_{-94}(2\sigma)\). We observe that the fitted value of \(M_h\) does not depend much on whether \(R_b\) is included or not into the fit. This is important in view of the large deviation in \(R_b\). However, some caution in the conclusions is still necessary: if both \(R_b\) and \(A_{LR}^e\) are absent from the fitted observables we get \(M_h = 205^{+226}_{-116}(1\sigma)^{+660}_{-170}(2\sigma)\). Thus, the data are consistent with a light Higgs boson but the \(2\sigma\) upper bound depends strongly on the inclusion of the SLD result for \(A_{LR}^e\) in the fit.

Another point of recent interest is the value of \(\alpha_s(M_Z)\) obtained from the electroweak fits We get \(\alpha_s(M_Z) = 0.122 \pm 0.005\) and this value is somewhat larger then the value obtained from the deep inelastic scattering data [2] \(\alpha_s(M_Z) = 0.112 \pm 0.005\).

We can interpret the SM fits as the MSSM fits with all superpartners heavy enough to be decoupled. Explicit calculations show that this occurs for superpartner masses already in the range 300-500 GeV. Supersymmetry then provides a rationale for a light Higgs boson: \(M_h \sim \mathcal{O}(100 \text{ GeV})\) and, since the best fit in the SM is consistent with the Higgs boson mass precisely in this range, the MSSM with heavy enough superpartners gives as good a fit to the precision electroweak data as the SM. This is the first interesting conclusion: the success of the SM is perfectly consistent with the supersymmetric Higgs sector and with the soft supersymmetry breaking scale in the range which maintains theoretical motivation for supersymmetry.

2. The \(R_b\) anomaly.

The MSSM with all superpartners heavy enough to decouple at the electroweak scale faces the same problem as the SM, namely of the \(R_b\) anomaly. The measurement of \(R_b\) is difficult and the possibility of larger than estimated systematic errors should be kept in mind. Still, the present result can at least be taken as a statistical hint towards a value of \(R_b\) somewhat larger than the SM prediction. It is then an interesting question if supersymmetry can enhance the value of \(R_b\). The issue has been addressed in a number of papers [3, 4, 5, 6, 7, 8, 9, 10]. It is well known already for some time that in the MSSM there are new contributions to the \(Z^0\bar{b}b\) vertex which can indeed significantly enhance the value of \(R_b\) (but do not change \(R_c\)) if some superpartners are sufficiently light [11, 3, 5, 6, 7, 8, 9]. More specifically, for low (large) \(\tan\beta\) the dominant contributions are chargino–stop (\(CP\)–odd Higgs boson and chargino–stop) loops. Moreover it is also known that new physics in \(\Gamma_{Z^0\to\bar{b}b}\) and therefore additional contribution to the total hadronic width of the \(Z^0\) boson would lower the fitted value of \(\alpha_s(M_Z)\) [12, 4, 13],
in better agreement with its determination from low energy data [2].

Any improvement in $R_b$ must maintain the perfect agreement of the SM with the other precision LEP measurements and must be consistent with several other experimental constraints (which will be listed later on). The bulk of the precision data, such as $M_W$, $\Gamma_Z$, $\sin^2 \theta_{\text{eff}}$, ..., are mostly sensitive to the $\Delta \rho$ parameter which measures the violation of the custodial $SU_V(2)$ symmetry. The contribution of the top–bottom quark mass splitting to $\Delta \rho$ leaves very little room for new contributions: $\Delta \rho < 0.0015$ at 95% C.L. [14]. Therefore, in order to maintain the overall good agreement of the fit with the data we must avoid new sources of the custodial $SU_V(2)$ symmetry breaking in the left currents. This is assured if the left squarks of the third generation and all left sleptons are sufficiently heavy [15, 6], say, $> O(300 \text{ GeV})$. At the same time, an increase in $R_b$ is sensitive mainly to the masses and couplings of the right handed top squark, charginos and - in the case of large $\tan \beta$ - of the right handed sbottom and $CP$–odd Higgs boson $A^0$ [11], which do not affect $\Delta \rho$ too much. Therefore, in the MSSM the requirement of a good overall fit is not in conflict with requirement of an enhancement of $R_b$ [6] and they imply a hierarchy:

$$M_{\tilde{t}_L} >> M_{\tilde{t}_R} \quad \text{or} \quad M_{\tilde{t}_2} >> M_{\tilde{t}_1}$$

(1)

with small left-right mixing.

The chargino - stop loop contribution to the $Z^0 b\bar{b}$ vertex can be with stop coupled to $Z^0$ and with charginos coupled to $Z^0$. In both cases the lighter the stop and chargino the larger is the positive contribution. Moreover, the $b\tilde{t}_i C^-$ coupling is enhanced for a right handed stop (it is then proportional to the top quark Yukawa coupling). Then, however, the stop coupling to $Z^0$ is suppressed (it is proportional to $g \sin^2 \theta_W$) and significant contribution can only come from the diagrams in which charginos are coupled to $Z^0$. Their actual magnitude depend on the interplay of the couplings in the $\tilde{C}_i^+ \tilde{t}_1 b$ and the $Z^0 \tilde{C}_i^+ C_j^-$ vertices. The first one is large only for charginos with large up-higgsino component, the second - for charginos with large gaugino component in at least one of its two-component spinors. It has been observed that, this combination never happens for $\mu > 0$. Large $R_b$ can then only be achieved at the expense of very light $C_j^-$ and $\tilde{t}_1$. In addition, for fixed $m_{C_1}$ and $M_{\tilde{t}_1}$, $R_b$ is larger for $r \equiv M_2/|\mu| > 1$ i.e. for higgsino-like chargino as the enhancement of the $C_i^- \tilde{t}_1 b$ coupling is more important than of the $Z^0 C_i^- C_j^-$ coupling.

For $\mu < 0$ the situation is different. In the range $r \approx 1 \pm 0.5$ a light chargino can be a strongly mixed state with a large up-higgsino and gaugino components (the higgsino-gaugino mixing comes from the chargino mass matrix). Large couplings in both vertices of the diagram with charginos coupled to $Z^0$ give significant increase in $R_b$ ($\sim 0.218$ for $m_{\tilde{t}} = 50 \text{ GeV}$)
even for the lighter chargino as heavy as $80 - 90 \text{ GeV}$ (similar increase in $R_b$ for $\mu > 0$ and the same stop masses requires $m_{C_1} \approx 50 \text{ GeV}$).

Significant enhancement of $R_b$ is also possible for large $\tan \beta$ values, $\tan \beta \approx m_t/m_b$ [11]. In this case, in addition to the stop–chargino contribution there can be even larger positive contribution from the $h^0$, $H^0$ and $A^0$ exchanges in the loops, provided those particles are sufficiently light (in this range of $\tan \beta$, $M_h \approx M_A$) and non-negligible sbottom–neutralino loop contributions. The main difference with the low $\tan \beta$ case is the independence of the results on the sign of $\mu$ (which can be traced back to the approximate symmetry of the chargino masses and mixings under $\mu \rightarrow -\mu$).

As stressed earlier an enhancement in $R_b$ must be subject to constraints from the quality of the global fit to the electroweak data and from all other available experimental information. Those constraints often differ in the degree of their model dependence and are worth careful discussion. A good quality of a global fit to the data is mainly assured by heavy enough left-handed sfermions with no direct impact on the value of $R_b$. The main remaining effect is the contribution of the decays $Z^0 \rightarrow N_i^0 N_j^0$ to the total and hadronic widths, $\Gamma_Z$ and $\Gamma_h$, respectively. One may wonder how this contribution depends on the assumption about the neutralino masses and their composition. With the GUT assumption, $M_1 = \frac{5}{3} \tan^2 \theta_W M_2$, the neutralino mass matrix is determined by the chargino one but, in general, one may consider arbitrary values of $M_1$. It is remarkable that for low $\tan \beta$ and for fixed ratio $r$ and $M_2$ the $\Delta \Gamma = \Sigma \Gamma_{Z \rightarrow N_i^0 N_j^0}$ is very weakly dependent on the value of $M_1$. In Fig.1a we show the contour $\Delta \Gamma = 5 \text{ MeV}$ for several values of $M_1$ and $\mu < 0$ (for $\mu > 0$ the $\Delta \Gamma$ is much smaller and is not relevant for the quality of the fit). The contributions of the individual channels change with a change of the neutralino masses and compositions but their sum remains almost constant when $M_1$ is changed arbitrarily. On the other hand, we see in Fig.1a that for fixed $m_C$ the dependence on the ratio $r$ is strong. For instance, for $m_C = 70 \text{ GeV}, \Delta \Gamma < 5 \text{ MeV}$ only for $r > 3$ or $r < 0.7$.

Similar information for large $\tan \beta$ is shown in Fig.1b. Here the results are symmetric with respect to the sign of $\mu$. The dependence of $\Delta \Gamma$ on the $M_1$ is not negligible and (this time for both signs of $\mu$) the larger the $M_1$ the larger the region in $(r, m_C)$ plane where $\Delta \Gamma < 5 \text{ MeV}$.

The $Z^0$ decay width into neutralinos has important impact on the quality of the global fit and constrains the region of light charginos. Thus, it puts a limit on the realistic increase in $R_b$. Fig.2 helps to understand this point in a clear way. We depict there the bound $\Delta \Gamma < 5 \text{ MeV}$ together with contours of constant $\delta R_b^{\text{SUSY}}$. The full results for $R_b$ as a function of $m_C$ (obtained from a scan over the parameters $\theta^L_t$, $\alpha_s(M_Z)$ with $M_{t_1} = 55$ GeV, $M_{t_2}$, $M_f$ (masses of other squarks and sleptons) fixed to 1 TeV; in the case of large $\tan \beta$, $M_A$ was fixed to 55 GeV and $M_{b_R} = 130$ GeV)
together with the corresponding values of the $\chi^2$ are shown in Fig.3a and 3b, for $\mu < 0$ and in Fig.4 for $\mu > 0$. The observed rise in the $\chi^2$ for light charginos is precisely due to large values of $\Delta\Gamma$.

The results presented in Fig.3b and 4 are valid as long as we do not assume the $R-$parity conservation $^1$. In this case the only additional constraints are the model independent bounds

$m_{C^-} > 47$ GeV,
$M_{\tilde{t}_1} > 46$ GeV,
$M_h > 60$ GeV,
$M_A > 55$ GeV (for large tan $\beta$).

The first two bounds are implicitly included in Fig. 3 and 4. The rôle of the lower limit on the Higgs boson mass (for a compact formula for radiatively corrected lighter Higgs boson mass in the limit $M_A >> M_Z$ see [16]) depends on the mass of the heavier stop and the left-right mixing angle. For $M_{\tilde{t}_2} > 500$ GeV (as required for good quality of the global fit) and small mixing angles (necessary for large $R_b$) $M_h$ is above the experimental limit $^2$ in a large range of the parameter space. Very small and large left-right mixing angles are, however, ruled out by this constraint $^3$. The Higgs boson mass bound has been imposed on the results of Fig. 3 and 4.

Assuming $R-$parity conservation, one has much stronger constraints which follow from the existence of the LSP and the corresponding decay signatures of the superpartners. The most important one are:

$m_{C^-} > 65$ GeV for $|m_{C^-} - m_{N^0}| > 5$ GeV $^4$,
$\Gamma(Z^0 \rightarrow N^0_1N^0_1) < 4$ MeV,
$BR(Z^0 \rightarrow N^0_1N^0_2) < 10^{-4}$,
D0 exclusion plot in $(M_{\tilde{t}_1},m_{N^0}^0)$ for $m_{C^-} > M_{\tilde{t}_1}$ $^5$,
$M_2 > 36$ GeV from the gluino search under the assumption $M_3 \approx 3M_2$ $^6$.

The bound on the $BR(Z^0 \rightarrow N^0_1N^0_2)$ is also shown in Fig.1 and 2. For small tan $\beta$ and $\mu < 0$ this bound is stronger than the effect of $\Delta\Gamma$ and puts a lower limit on $m_{C^-}$ which is often stronger than the direct limit of 65 GeV (see Fig.3c where the full results obtained under the assumption of $R-$parity conservation are presented). The values of the $\chi^2$ as a function of the chargino mass (in the now allowed range of $m_{C^-}$) are almost identical to Fig.3a.

For low tan $\beta$ and $\mu > 0$ the constraint disappears and only the direct

$^1$The direct contribution to $R_b$ from the diagrams with the $R-$parity violating couplings is generically negligible [27] and is not included here. We merely discuss the rôle of the assumption about $R-$parity conservation in constraining the dominant MSSM contribution to $R_b$.

$^2$Important rôle of the experimental lower bound on $M_h$ in ref. [9] in constraining the potential increase of $R_b$ is due to the chosen upper bound $M_{\tilde{t}_1} < 250$ GeV which, anyway, looks too low from the point of view of a global fit.
limit on \( m_C \) is relevant. For large \( \tan\beta \) (and both signs of \( \mu \)) the bound is relevant only for \( r < 1 \) whereas the increase in \( R_b \) comes mainly for \( r > 1 \). Therefore, the results shown in Fig.4 do not depend on whether \( R^- \) parity is conserved or explicitly broken.

Light chargino and stop as well as light charged Higgs boson induce non-standard top quark decays (\( t \rightarrow \tilde{t}_1 N_i^0 \) and \( t \rightarrow H^+ b \)). Present experimental limits on those branching ratios are very uncertain. On the theoretical side, e.g. the branching ratio for \( t \rightarrow \tilde{t}_1 N_i^0 \) in the low tan \( \beta \) region depends strongly on the properties of the neutralinos and, generically, is acceptable for a gaugino - like neutralino [19]. In Fig.5 we show contours of this branching ratio for \( \tan\beta = 1.4 \) and \( \mu < 0 \), with the other parameters fixed to the values which give the results for \( \delta R_b \) which are shown in Fig.2.

Finally, light chargino and stop play important role in the FCNC transitions. This is discussed in more detail in the next section. This last constraint is also included in the full results for \( R_b \) presented in Fig.3 and 4.

The overall summary of the results presented in Fig.3 and 4 is as follows: Even with no assumption about the \( R^- \) parity conservation and for arbitrary \( M_1 \), the realistic upper bound for \( R_b \) in the MSSM is \( \sim 0.218 \) (0.2185) for low (large) \( \tan\beta \) values and for both signs of \( \mu \). Further increase is bounded by too large values of \( \Delta \Gamma \). In this general case, and with maximal \( R_b \), the chargino mass is as low as 50 GeV for small \( \tan\beta \) and \( \mu > 0 \) but \( m_C > 60 \) GeV for \( \mu < 0 \) and always for large \( \tan\beta \). In the MSSM with \( R^- \) parity conservation, the increase in \( R_b \) is even more strongly constrained (mainly by the direct limit \( m_C > 65 \) GeV and by the \( BR(Z \rightarrow N_1^0 N_2^0) \)) and generically smaller by \( \sim 0.001 \), except for the region of low \( \tan\beta \) (\( \sim 1 - 2 \)) and \( \mu < 0 \). There, \( R_b \sim 0.218 \) is still realistic for \( m_C \sim (70 - 90) \) GeV. The right-handed stop can be around 50 GeV but even with \( M_1 \sim (60 - 70) \) GeV the effect on \( R_b \) is not negligible. An increase in \( R_b \) gives a shift in the fitted value of \( \alpha_s(M_Z) \): \( \delta \alpha_s = -4 \delta R_b \). Therefore, \( \delta R_b \sim (0.001 - 0.002) \) gives \( \alpha_s(M_Z) \sim (0.118 - 0.114) \). Finally, it is important to observe that the maximal enhancement in \( R_b \) in the MSSM compared to the SM value is only of the order of the 1.5\( \sigma \) present experimental error. Significant improvement of the experimental precision is necessary to confirm such a difference between the two theories.

### 3. Rare processes with light chargino and stop.

There are several well known supersymmetric contributions to rare processes. In particular, supersymmetry may provide new sources of flavour violation in the soft terms. However, even assuming the absence of such new effects, there are obvious new contributions when \( W^\pm - q \) SM loops are replaced by the \( H^\pm - q \) loops and by \( \tilde{W}^\pm (\tilde{H}^\pm) - \tilde{q} \) loops. Those can be expected to be very important in the presence of light chargino and stop.
and they contribute to all best measured observables: $\epsilon$ parameter for the $K^0 - \bar{K}^0$ system, $\Delta m_B$ from $B^0 - \bar{B}^0$ mixing and $BR(b \to s\gamma)$.

There are two important facts to be remembered about these contributions. They are present even if quark and squark mass matrices are diagonal in the same super-Kobayashi-Maskawa basis. However the couplings in the vertex $d_i \bar{u}_j C^-$ depend on this assumption and can depart from the $K$-$M$ parametrization if squark mass matrices have flavour-off diagonal entries in the super-Kobayashi-Maskawa basis. Some of those entries are still totally unconstrained and this is precisely the case for the (right) up squark mass matrix which is relevant e.g. for the couplings $b \bar{t}_R C^-$. Still, sizeable suppression compared to the $K$-$M$ parametrization requires large flavour-off diagonal mass terms, of the order of the diagonal ones. To remain on the conservative side we include the constraints from rare processes under the assumption of the $K$-$M$ parametrization of the chargino vertices. The rôle of the $b \to s\gamma$ constraint is discussed in detail in ref. [10]. The requirement of the acceptable $BR(b \to s\gamma)$ has been imposed on the results shown in Fig.3 and 4.

The second important remark is that the element $V_{td} \approx A\lambda^3(\rho - i\eta)$ (in the Wolfenstein parametrization), which is necessary for the calculation of the chargino and charged Higgs boson loop contribution to the $\epsilon_K$ parameter and the $B^0 - \bar{B}^0$ mixing, is not directly measured. Its SM value can change after the inclusion of new contributions. Thus the correct approach is the following one: take e.g.

$$\Delta m_B \approx f_{B_d}^2 B_{B_d} |V_{tb}V_{td}^*|^2 |\Delta|$$

where

$$\Delta = \Delta_W + \Delta_{NEW}$$

is the sum of all box diagram contributions, $f_{B_d}$ and $B_{B_d}$ are the $B^0$ meson decay constant and the vacuum saturation parameter. The $CP$ violating parameter $\epsilon_K$ can also be expressed in terms of $\Delta$. Given $|V_{cb}|$, and $|V_{ub}/V_{cb}|$ (known from the the tree level processes i.e. almost unaffected by the supersymmetric contributions) one can fit the parameters $\rho$, $\eta$ and $\Delta$ to the experimental values of $\Delta m_{B_d}$ and $|\epsilon_K|$. This way we find [20, 21] a model independent constraint

$$\frac{\Delta}{\Delta_W} < 3$$

for $\sqrt{f_{B_d}^2 B_{B_d}}$ in the range $(160 - 240)$ GeV and $B_K$ in the range $(0.6 - 0.9)$ GeV, which in the next step can be used to limit the allowed range of the stop and chargino masses and mixings. The parameter space
which is relevant for an increase in $R_b$ gives large contribution to $\Delta$. It is still consistent with the bound (3) but requires modified (compared to the SM) values of the $CP$–violating phase $\delta (\eta, \rho)$.

4. $R_b$ and exotic events.

A single event $e^+ e^- \gamma \gamma + \text{missing } E_T$ has been reported by Fermilab [22]. Preliminary results from the LEP1.5 run (after the upgrade of energy to $\sqrt{s} = 130 - 136$ GeV) include peculiar four-jet events reported by ALEPH [23]. Although statistics is too low to exclude fluctuations, it is interesting to speculate if they can be explained by supersymmetry and whether simultaneous explanation of these events and the $R_b$ anomaly is possible. A detailed study of the Fermilab event in the supersymmetric extension of the SM is a subject of refs. [24]. It is interesting to observe that the Fermilab event can be explained as a selectron pair production, with the supersymmetric spectrum which is consistent with larger than in the SM values of $R_b$. The best description is obtained for $M_1 \approx M_2$ but in a model with $R$–parity conservation. This last fact should be stressed in view of the following discussion of ALEPH events.

ALEPH 4-jet events have very peculiar gross features. On the kinematical grounds they can be interpreted as a production of a pair of new particles $X$ with $m_X \approx 55$ GeV and a relatively large effective (after cuts) production cross section $\sigma \approx 3.7 \pm 1.7$ pb. Any interpretation of this new particle is strongly constrained by the decay signature: to a good approximation no missing energy has been observed and most of the events do not contain identified $b$–quark jets and no fast leptons in the final state. Those signatures of the events imply that any explanation within a $R$–parity conserving MSSM is very difficult (for an explanation based on the idea of a light gluino, $m_{\tilde{g}} \sim 1$ GeV see ref. [25]). Moreover, a large production cross section is not easy to accommodate ($R$–parity violation has little impact on the production cross section so it can be reliably estimated in the MSSM). A sneutrino pair production seems to be an acceptable possibility [26] but its connection to the $R_b$ anomaly is not obvious. Turning now to supersymmetric fermions, a neutralino of a mass 55 GeV has production cross section more than one order of magnitude below the reported value. Thus, we are left with a light chargino as the most interesting candidate to explain ALEPH events. Indeed, the full production cross section are typically large $\mathcal{O}(10 \text{pb})$ (see Fig.6). Moreover, there is an interesting link with $R_b$ anomaly.

The question which remains is whether chargino decay signatures can be consistent with ALEPH data. No missing energy rules out $R$–parity conserving schemes. If $R$–parity is not conserved, additional terms in the
superpotential are allowed

\[ W = \frac{1}{2} \lambda_{ijk} L_i L_j E^c_k + \lambda'_{ijk} L_i Q_j D^c_k + \frac{1}{2} \lambda''_{ijk} U^c_i D^c_j D^c_k \]  \hspace{1cm} (5)

where \( \lambda_{ijk} = -\lambda_{jik} \) and \( \lambda''_{ijk} = -\lambda''_{ikj} \). The first two terms violate lepton number conservation and the last one - baryon number conservation. Simultaneous presence of both types of terms can lead to a rapid proton decay. However, only \( \lambda \) and \( \lambda' \) or only \( \lambda'' \)-type couplings are allowed, of course within the present experimental limits. The latter depend on the type of the coupling but for several of them are relatively weak, particularly for the couplings involving the third family.

If the chargino decayed through a lepton number violating coupling, there should be a hard lepton or missing momentum in the event. Thus this explanation looks unlikely. With baryon number violating couplings \( \lambda'' \), the chargino may decay via either of two channels

\[ C^\pm \rightarrow \tilde{q}_1^* q_2 \rightarrow q_2 q_3 q_4 \] \hspace{1cm} (6)

where the squark is right handed and can be virtual, and

\[ C^\pm \rightarrow N^0* W^\pm* \rightarrow N^0* f_1 f_2 \rightarrow q_3 q_4 q_5 \] \hspace{1cm} (7)

with \( N^0 \) real or virtual. The actual decay pattern depends on the details of the couplings and the values of the masses. It was shown in ref. [27] that, with most natural assumptions, that right stop is the lightest squark and that the coupling \( \lambda''_{tds} \) is the largest one, the decay mode (6) is the more interesting one. Although at the qualitative level we may expect in this scenario too many jets and/or \( b \)-quark jets in the final state, it happens that the effective 4-jet cross section (after the experimental cuts) with at most one \( b \)-quark contributing to the visible energy is of the right order of magnitude (see Fig.7a). Particularly attractive is the possibility (Fig.7b) \( m_{C^-} \gtrsim M_{\tilde{t}_R} \approx 55 \text{ GeV} \) and with the mass difference such that the stop is real but produced almost at rest. Then ALEPH events can be explained by [27]

\[ Z^0 \rightarrow C^- C^+ \rightarrow (\tilde{t}_R \bar{b})(\tilde{t}_R b) \rightarrow (\bar{d}s\bar{b})(dsb) \] \hspace{1cm} (8)

with very slow \( b \)-quarks and therefore escaping detection. With the present experimental resolution, a mass degeneracy \( m_{C^-} - M_{\tilde{t}_R} \lesssim 5 - 12 \text{ GeV} \) is sufficient for this scenario [27]. This case is unique in the sense of avoiding
b-jets as well as giving the correct jet invariant mass distribution with about half of the events falling into the range

\[ 103.3 \text{GeV} < \Sigma m(= m_{ij} + m_{kl}) < 106.6 \text{GeV} \quad (9) \]

The results for the chargino decay via virtual stop shown in Fig.6a generically give much broader distributions for this sum of the properly paired two jet invariant masses.

Neutralino could still be light but the decay \( C^- \rightarrow N^0 f_1 f_2 \) is suppressed due to kinematical reasons (due to multibody final states). A link with the \( R_b \) anomaly is clear. However, simultaneous supersymmetric explanation of the Fermilab (one) event and the ALEPH 4−jet events looks unlikely because of the need for broken \( R \)−parity in the latter case.

5. Conclusions.

Chargino and right handed stop are likely to be (in addition to neutralinos) the lightest supersymmetric particles. Not only the masses in the range or even below \( M_Z \) are not excluded by any presently available experimental data, they may be responsible for an increase in \( R_b \) up to \( \sim 0.218 \). Depending on whether \( R \)−parity is conserved or not, the Fermilab event or ALEPH events may be explained simultaneously with larger than the SM values of \( R_b \). In particular, ALEPH events, if confirmed, may signal the discovery of a chargino with \( m_C \approx 60 \text{ GeV} \), a stop with \( M_{\tilde{t}} \approx 55 \text{ GeV} \) and broken \( R \)−parity.

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see also S. Bethke talk at Rencontres de Moriond, March 1995.

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Figure 1: Dependence of the contour $\Delta \Gamma \equiv \Sigma \Gamma_{Z \rightarrow N^0_i N^0_j} = 5$ MeV on $M_1$. $M_1 = 0.2M_2$ (dashed), $M_1 = 0.5M_2$ (solid) and $M_1 = M_2$ (dotted). Dash-dotted lines show $BR(Z \rightarrow N^0_i N^0_j) = 10^{-4}$ for $M_1 = 0.2M_2$ (upper ones) and $M_1 = 0.5M_2$ (lower ones). In a), the region below the lower solid line is excluded by “kinematics” (the structure of the chargino mass matrix).
Figure 2: Contours of $\delta R_b^{SUSY} = 0.0020, 0.0015$ and $0.0010$ (dash-dotted lines) in the plane ($M_2/\mu, m_C$) for $m_t = 170$ GeV and different $\tan \beta$. Also shown are the lines of $\Delta \Gamma = 5$ MeV (dashed) and $BR(Z^0 \to N_1 N_2) = 10^{-4}$ (dotted). Solid lines show the “kinematic” limits of $m_C$ and $M_\tilde{g} = 150$ GeV.
Figure 3: Best $\chi^2$ and the corresponding $R_b$ as a function of $m_C$, for negative $\mu$, $m_t = 170$ GeV, $M_{\tilde{t}_1} = 50$ GeV, $M_{\tilde{t}_2} = M_A = 1$ TeV and for two values of $\tan \beta$. For $R_b$ two cases are shown: with $R$–parity broken and with $R$–parity conserved. $\chi^2$ as a function of $m_C$ is the same in both cases. Different lines correspond to different values of $M_t^2/\mu$: 0.2 (upper (lower) solid for a (b,c)), 0.5 (dashed), 1 (dotted), 1.5 (dash-dotted) and 3 (lower (upper) solid for a (b,c)).
Figure 4: Best $\chi^2$ and the corresponding $R_b$ as a function of $m_C$, for positive $\mu$, $m_t = 170$ GeV, $M_{\tilde{t}_1} = 50$ GeV, $M_{\tilde{t}_2} = 1$ TeV, $M_A = 1$ TeV for $\tan \beta = 1.4$ and $M_A = 55$ GeV for $\tan \beta = 50$. Different lines correspond to different values of $M_2/\mu$. For $\tan \beta = 1.4$: 0.5 (solid), 1 (dashed), 1.5 (dotted) and 3 (dash-dotted); for $\tan \beta = 50$: 1 (upper (lower) solid for $\chi (R−b)$), 1.5 (dashed), 3 (dotted), 5 (dash-dotted) and 10 (lower (upper) solid for $\chi (R_b)$). Identical results are obtained for broken or conserved $R$–parity.
Figure 5: Contours of $BR(t \to \tilde{t}_1 N^0) = 25\%$ and $45\%$ for $\tan \beta = 1.4$, $\mu < 0$ and the other parameters fixed at the values which give $\delta R_b$ as shown in Fig.2. The solid curve is the “kinematical” limit for the chargino mass.
Figure 6: Cross sections for 55 GeV chargino production for different choices of \((\tan\beta, M_{\tilde{\nu}})\) values: for \(\mu > 0\): (1.4, 50)-solid, (1.4, 200)-dashed, (50, 50)-dotted, (50, 200)-dashdotted; for \(\mu < 0\): (1.85, 50)-solid, (1.85, 200)-dashed, (50, 50)-dotted, (50, 200)-dashdotted.
Figure 7: Effective 4-jet cross sections for pair-produced charginos decaying through the stop channel: a) for 55 GeV chargino, requiring that at most one b-quark contributes to the visible energy, b) for 59 GeV chargino decaying through real 52 GeV right stop (the lower (upper) curves are with (without) the cut of eq.9).