Seeking supersymmetry at LEP

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• Scope and motivation

This minireview has been commissioned by the organizers of the workshop. Under the general title of LEP, I aim to cover three segments of the experimental programme in that accelerator. First, I consider the analyzed sample of $Z$-events out of the 20 million collected at LEP 1. The lack of observation of any direct supersymmetry signal in these has yielded bounds on supersymmetric parameters. Second, there is about 5 (pb)$^{-1}$ of data which have been collected during the intermediate energy run (LEP 1.5) at 130, 136 and 140 GeV in the $e^+e^-$ centre of mass. A preliminary examination of these data is revealing for supersymmetry search strategies. Finally, I briefly indulge in some futurology vis-a-vis the coming runs in LEP 2 which will be in the energy regime 161-200 GeV.

Our supersymmetry reference point is the supergravity-constrained [1] minimal supersymmetric standard model CMSSM – distinguished by the possession of the smallest number of parameters among all contending SUSY scenarios. Sometimes, a less strong version of this model is considered [2] in which one partially relaxes the supergravity-based unification assumption of a universal supersymmetry-breaking scalar mass at the scale by taking separate scalar mass values for the Higgs and the sfermion sectors; this version will be differentiated by using the prefix “partially constrained”: PCMSSM. Once in a while, we shall also mention the supergravity-constrained next-to-minimal supersymmetric standard model CNMSSM [3], possessing an additional gauge-singlet superfield $N$ whose scalar VEV $\langle n \rangle$ leads to the Higgsino mass parameter $\mu$.

The motivation for considering sub-TeV to TeV scale supersymmetry in particle physics is well-known. It naturally ensures the stability of the weak scale against radiative perturbations induced from unknown high scale physics. It has a pronounced decoupling feature in that at low energies it reduces to the Standard Model with a somewhat light physical Higgs scalar when all sparticles are made heavier [1] than about 200 GeV. As a result, all the presently successful features of the Standard Model are easily retained (in the sense of being compatible within errors) in the supersymmetric extension. A particularly notable item among the latter is quantum consistency at the 1-loop level which predicts [5] the top quark mass $m_t$ to be $169 + 7^{+4}_{-3}$ GeV to be compared with the Tevatron-measured [summer, 1996] value of $175 \pm 9$ GeV from CDF and $170 \pm 18$ GeV from D0. In addition, though, supersymmetry provides the standard model with a bridge from electroweak energies to very high scale physics, e.g. $M_X \sim 3 \times 10^{16}$ GeV [6] of supersymmetric grand unification or $M_S \sim 5 \times 10^{17}$ GeV of superstring unification [6] or the Planck scale $M_P \sim 1.2 \times 10^{19}$ GeV of quantum gravity.

Apart from the above virtues, the CMSSM is characterized by a powerful simplicity as well as phenomenological tractability and definite testability. As many as 31 parameters, which appear [8] in the straightforward minimal extension of the low-energy Standard Model to its supersymmetric version (MSSM) with soft supersymmetry breaking (SSB)

$$L^{SM} \rightarrow L^{SSM} + L^{SSB},$$  \hspace{1cm} (1)

are reduced in the supergravity-constrained MSSM to 4 plus a sign – namely that of $\mu$. (In PCMSSM one has 5 parameters including the magnitude of $\mu$). Moreover, there are well-formulated tests, the violation of any of which will kill the model. As an illustration, one may
mention the possible observation of a slepton heavier than the corresponding squark (in the first two generations) which will destroy CMSSM. Another killer would be the demonstration of the absence of any neutral Higgs particle with a mass below about 130 GeV. The latter, in fact, would exclude MSSM itself.

Let us focus on the crucial assumptions that go into the CMSSM chosen here and its salient features. To start with, there is a spectrum comprising particles and their partner sparticles. The particles are those of the Standard Model except that, instead of one, there are five physical Higgs scalars (two charged $H^\pm$, two CP-even neutrals – the lighter $h$ and the heavier $H$ – and one CP-odd neutral $A$) coming from two Higgs doublets. (The ratio of the VEV of the neutral Higgs field which couples with up-type fermions to that of the one which does so with down-type fermions – is called $\tan \beta$). The generic sparticle is expected to be heavier than the corresponding particle, though the order could be reversed for the right-chiral stop and/or the chargino, neutralino. All coupling and mass parameters evolve with the energy scale – described by 26 independent renormalization group equations. Next, postulated boundary conditions are imposed on some of these parameters at the unification scale $M_X \sim 2 \times 10^{16}$ GeV. Specifically, all supersymmetry-breaking scalar (gaugino) masses are taken to be universal and equal to one mass $m_0$ ($M_{1/2}$). Squared masses of the Higgs at the unification scale of course, have an additional supersymmetric contribution, namely $\mu^2$, since the higgsino mass is also a supersymmetric contribution to the Higgs mass. Again, all supersymmetry breaking trilinear scalar couplings, $A_{ijk}$ ($i, j, k$ generation indices) are taken to be equal ($= A_0$) at $M_X$. Here $m_0$ and $M_{1/2}$ are supposedly of the order of the gravitino mass $m_{3/2} \sim \Lambda^2_S M_{pla}^{-1}$, with $\Lambda_S$ being some dynamical supersymmetry-breaking scale in the hidden sector $\sim 10^{11} - 10^{12}$ GeV. Now $m_0, A_0, m_{1/2}$ and $\tan \beta$ can be chosen to be the four parameters; or, $m_A$ could be traded for one of the first two. The Higgs squared masses evolve rapidly in a decreasing manner as one lowers the energy scale. The starting values of the parameters at the unification scale ensure that one such squared mass of a neutral Higgs flips its sign, while evolving, thus leading to radiative electroweak symmetry breakdown and a weak scale $M_W < O(m_{3/2})$. Though CMSSM still has a large allowed region in its space of parameters, we will – from time to time – talk about two extreme scenarios just to emphasize their difference:

# Decoupling SUSY — all sparticles heavier than 200 GeV.

# SUSY round the corner — a stop, a chargino + a neutralino, each below 100 GeV.

In the less ambitious PCSSM, universality at $M_X$ is imposed only on squark and slepton (and not on the Higgs) masses – thus keeping an open mind on the radiative origin of the Higgs mechanism. Caveat emptor: Some of the highly restrictive assumptions underlying CMSSM may be plain wrong!

- Lightning review of CMSSM

The superfield content of the model in the matter sector, written in a transparent notation ($i = 1, 2, 3$ is a generation index), is:

$$Q_L^i = \left( \begin{array}{c} U_L^i \\ D_L^i \end{array} \right), \quad L_L^i = \left( \begin{array}{c} N_L^i \\ E_L^i \end{array} \right), \quad Q_R^i, L_R^i, H_1 = \left( \begin{array}{c} H_1^0 \\ H_1^- \end{array} \right), \quad H_2 = \left( \begin{array}{c} H_2^+ \\ H_2^0 \end{array} \right).$$
The corresponding superpotential (with R-parity assumed conserved) is:

$$ W = \lambda^U_i Q^i_L H_2 U^j_R + \lambda^D_i Q^i_L H_1 D^j_R + \lambda^E_i L^i_L H_1 E^j_R + \mu H_1 H_2, $$

with \( \lambda \)'s as Yukawa couplings.

The scalar potential can be derived from (1). Writing \( \phi_j \) for a generic scalar field and incorporating the soft supersymmetry breaking terms, we have

$$ V = \sum_j \left| \frac{\partial W}{\partial \phi_j} \right|^2 + D\text{-terms} + \sum_{i,j} m^2_{ij} \phi_i \phi_j $$

$$ + \left\{ A_U \lambda_U q^L h^2 u_R + A_D \lambda_D q^L h^1 d_R + A_E \lambda_E e^L h^1 e_R + B \mu h_1 h_2 + H.C. \right\}. $$

In (3) the third RHS term includes \( m_1^2 h^1_R h^1_L + m_2^2 h^2_R h^2_L \) with a vanishing \( m_{1,2} \) where \( h^1_{1,2} \) refers to the scalar component of the superfield \( H_{1,2} \). Also, \( v_{1,2} = \langle h^0_{1,2} \rangle \) and \( \tan \beta = v_2/v_1 \). The physical fields can be expressed in terms of the superfield components given above. For instance, the field for the lightest neutral scalar is

$$ h = \sqrt{2}(\text{Re} h^0) \cos \alpha - \sqrt{2}(\text{Re} h^0 - v_1) \sin \alpha, $$

where \( \alpha \) is an angle which enters via mixing. The orthogonal heavier combination is

$$ H = \sqrt{2}(\text{Re} h^0 - v_2) \sin \alpha + \sqrt{2}(\text{Re} h^0 - v_1) \cos \alpha \quad \text{while} \quad A \text{ equals } \sqrt{2}(\text{Im} h^0 \cos \beta - \text{Im} h^0 \sin \beta). $$

The partners of the CKM matrices in the scalar sector are assumed to possess safety properties which suppress dangerous flavor-changing neutral current processes that could emerge from (1).

At the tree level itself one has several mass relations.

$$ m^2_{\pm} = m^2_A + M^2_W, $$

$$ m^2_h \leq M^2_Z \leq m^2_H, $$

$$ \frac{m_h}{\cos 2\beta} < m_A < m_H, $$

$$ \mu^2 = (\cos 2\beta)^{-1}(m^2_2 \sin^2 \beta - m^2_1 \cos^2 \beta) - \frac{1}{2} M^2_Z, $$

$$ 2B\mu = (m^2_1 - m^2_2) \tan 2\beta + M^2_Z \sin 2\beta, $$

$$ m^2_A = m^2_1 + m^2_2 + 2\mu^2. $$

On including 1-loop quantum corrections in the leading log approximation, the upper bound on the squared mass of \( h \) reads \( (\tilde{t}_{1,2} \text{ are the two physical squarks, assumed to weigh more than the top) [4]:} \)

$$ M^2_h < M^2_Z \cos^2 2\beta + \frac{3\alpha_{EM}}{2\pi \sin^2 \theta_W} \frac{m^4_t}{M^2_W} \ln \frac{m_t m_{\tilde{t}_1} m_{\tilde{t}_2}}{m^2_t} \simeq (130 \text{ GeV})^2. $$

The boundary conditions at \( M_X \) imply

$$ m^2_1(M_X) = m^2_2(M_X) = m^2_0. $$
Turning to gaugino masses $M_i$ ($i = \text{nonabelian gauge group index}$) and considering 1-loop RGE effects, one can write – with $\alpha_u$ as the unified fine structure coupling –

$$M_i(Q) = M_{1/2} \alpha_i(Q) \alpha_u^{-1}(M_X).$$  \hspace{1cm} (7)

For the $U(1)_Y$ case, with the standard definition of $Y$, there is an extra factor of $5/3$ in the RHS. It turns out that $M_1$ ($M_Z$) $\simeq 0.41$ $M_{1/2}$ and $M_2$ ($M_Z$) $\simeq 0.84$ $M_{1/2}$ with a mild $Q$-dependence in $M_{1/2}$. However, the situation is quite different for $M_3$. The physical on-shell gluino mass $m_{\tilde{g}}$ is given by \[1\]

$$m_{\tilde{g}} = M_3(Q) \left[ 1 + \frac{\alpha_s(Q)}{4\pi} \left\{ 15 - 18 \ell_n \frac{M_3(Q)}{Q} + \sum_i \int_0^1 dx \frac{x \ln x}{m_i^2 + (1-x)m_{\tilde{g}}^2 - x(1-x)M_3^2} \right\} \right]$$ \hspace{1cm} (8)

and is independent of $Q$. For $M_3 \simeq 0.1 \text{ TeV}$ and $m_{\tilde{g}} \simeq 1 \text{ TeV}$, the difference between $m_{\tilde{g}}$ and $M_3$ ($M_3$) can be as much as 30%.

The spectrum of the remaining sparticles can be parametrized, after accounting for renormalization group evolution, as follows \[11\]:

$$m_{\tilde{e}_R}^2 = m_0^2 + 0.15 M_{1/2}^2 - \sin^2 \theta_W D;$$  \hspace{1cm} (9a)

$$m_{\tilde{e}_L}^2 = m_0^2 + 0.52 M_{1/2}^2 - \left( \frac{1}{2} - \sin^2 \theta_W \right) D;$$  \hspace{1cm} (9b)

$$m_{\tilde{g}}^2 = m_0^2 + 0.52 M_{1/2}^2 + \frac{1}{2} D;$$  \hspace{1cm} (9c)

$$m_{\tilde{q}_1 R}^2 = m_0^2 + (0.07 + C_g) M_{1/2}^2 + \frac{2}{3} \sin^2 \theta_W D;$$  \hspace{1cm} (9d)

$$m_{\tilde{q}_1 L}^2 = m_0^2 + (0.07 + C_g) M_{1/2}^2 - \frac{1}{3} \sin^2 \theta_W D;$$  \hspace{1cm} (9e)

$$m_{\tilde{q}_2 L}^2 = m_0^2 + (0.47 + C_g) M_{1/2}^2 + \left( \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right) D;$$  \hspace{1cm} (9f)

$$m_{\tilde{q}_2 L}^2 = m_0^2 + (0.47 + C_g) M_{1/2}^2 - \left( \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right) D;$$  \hspace{1cm} (9g)

Here $C_g = \frac{8}{9} [\alpha_S^2(m_{\tilde{g}})/\alpha_S^2(M_X) - 1]$ and $D = M_2^2 \cos^2 \beta$ while we have $\ell = e, \mu, q^+ = u, c$ and $q^- = d, s, b$. For stops and staus, considerable left-right mixing is anticipated. The corresponding mass-squared matrices are given by

$$m_{\tilde{t}}^2 = \begin{pmatrix} m_{\tilde{t}_L}^2 + m_{\tilde{t}}^2 + 0.35 D & -m_{\tilde{t}}(A_t + \mu \cot \beta) \\ -m_{\tilde{t}}(A_t + \mu \cot \beta) & m_{\tilde{q}_1 L}^2 + m_{\tilde{t}}^2 + 0.16 D \end{pmatrix};$$  \hspace{1cm} (10a)

$$m_{\tilde{\tau}}^2 = \begin{pmatrix} m_{\tilde{\tau}_L}^2 + m_{\tilde{\tau}}^2 - 0.27 D & -m_{\tilde{\tau}}(A_\tau + \mu \tan \beta) \\ -m_{\tilde{\tau}}(A_\tau + \mu \tan \beta) & m_{\tilde{\tau}_R}^2 + m_{\tilde{\tau}}^2 - 0.23 D \end{pmatrix};$$  \hspace{1cm} (10b)

A sample scatter plot of the ranges \[12\] of some characteristic masses in the model – showing the extent of variation in the parameter space – is shown in Fig. 1. One should also mention
that five squarks (i.e. all except the stop) need to be taken as nearly mass-degenerate in order to avoid an unacceptable FCNC-induced $K^0 - \bar{K}^0$ mixing. This could be a problem in Fig. 1 [12] which has a rather large $\bar{b}_L - \bar{b}_R$ mass-splitting. A similar argument vis-a-vis the FCNC-induced $\mu \rightarrow e\gamma$ decay requires the near mass-degeneracy of all sleptons except $\tilde{\tau}$.

- **LEP 1 bounds and LEP 1.5 results**

  **Higgs constraints**

  Let me first reiterate the experimental mass bounds on the neutral Higgs scalars. The SM Higgs particle has been searched for in the Bjorken process $e^+e^- \rightarrow Z \rightarrow hZ^* \rightarrow b\bar{b}\ell\bar{\ell}(q\bar{q})$. The present SM lower bound of 63.5 GeV gets diluted significantly in case of $h$ of MSSM since $\Gamma_{Bj}^{MSSM} = \sin^2(\alpha - \beta)\Gamma_{Bj}^{SM}$ and the situation has not been helped by the LEP 1.5 data. The exact lower bound on $m_h$ in MSSM is presently a matter of controversy. The CP-odd scalar $A$ has also been looked for in the process $e^+e^- \rightarrow Z \rightarrow hA$. By use of the formula

  $$\Gamma(Z \rightarrow hA) = \frac{1}{2} \cos^2(\alpha - \beta)\Gamma(Z \rightarrow \nu\bar{\nu}) \left[\lambda \left(1, m_h^2 M_Z^{-2}, m_A^2 M_Z^{-2}\right)^3\right]^{3/2}, \quad (11)$$

  the lack of observation of any $A$ translates to the bound $m_A > 26$ GeV for $1 < \tan\beta < 50$, assuming that $h$ allows the process kinematically. It is worth pointing out in this context that no MSSM Higgs production interpretation can be given to the excess $4j$ events reportedly seen by ALEPH at LEP 1.5 owing to the alleged lack of $b$’s in the final state.
Direct sparticle mass bounds

Charginos and charged sfermions are open to pair-production from the resonant $Z$. The signals would be hard acollinear jets (and or leptons) and $E_T$. The lack of observation of any such event at LEP 1 implies that those sparticles are all heavier than $M_Z/2$. In LEP 1.5, considering the offshell $Z$, the relevant lower bound would be $\sqrt{s}/2$. One also needs to take into account the experimental constraint that any possible partial width of the $Z$ decaying into visible channels must obey $\Delta \Gamma_Z$ (new and visible) < 0.13 MeV. All these constraints, combined within PCMSSM and assuming two light neutralinos, exclude certain regions of the $\mu - \tan \beta$ plane. We show these exclusion zones [13] in Fig. 2. Fig. 2a shows the situation for the decoupling case where all superparticles, except the lightest chargino, weigh more than 200 GeV; the boundaries of the dashed region, already excluded by LEP 1 data, would expand further as shown in case there is no chargino upto 90 GeV. A similar plot is made for the “round the corner” scenario in Fig. 2b.

![Figure 2: PCMSSM exclusion regions in the $\mu - \tan \beta$ plane](image)

Let me come next to neutral sparticles, such as sneutrinos and neutralinos. One needs to first consider the nonaccelerator constraints on the former. Stable sneutrinos in the range 3 GeV to 1 TeV are practically excluded by dark-matter search experiments. Unstable ones must weigh more than 41.8 GeV, otherwise they would allow the decay chain $Z \rightarrow \tilde{\nu}\tilde{\nu}^*$, $\tilde{\nu} \rightarrow \nu + \text{LSP}$ violating the LEP bound $\Delta \Gamma_Z$ (new and invisible) < 6.7 MeV. Turning to the neutralinos ($\tilde{\chi}_{1,2,3,4}$), we note that, among the modes $Z \rightarrow \tilde{\chi}_i^0 \tilde{\chi}_j^0$, the $\tilde{\chi}_1^0 \tilde{\chi}_1^0$ final state is undetectable and that the sum of the branching ratios of all such detectable processes is experimentally constrained to be less than $10^{-5}$. The simplest produced final state is $\tilde{\chi}_2^0 \tilde{\chi}_2^0$ with the subsequent decay $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 + \text{visible}$. There can, however, be other states [14] such as $\tilde{\chi}_3^0 \tilde{\chi}_4^0$, each of which can decay into $\tilde{\chi}_1^0 + Z$ or $h$. Analyzing the LEP data, one can conservatively quote $M_{\tilde{\chi}_1^0} > 23$ GeV in the CMSSM and PMSSM, but the lower bound reduces to 0 in CNMSSM. For $\tilde{\chi}_2^0$, a stronger statement, namely $M_{\tilde{\chi}_2^0} > 56$ GeV is possible [14].
Indirect constraints

These are obtained by inputting $\Gamma_{\text{tot}}, \Gamma_{\text{ee}}, \sin^2 \theta_{\text{EFF}}, A_{FB}^b, A_{FB}^c, \Delta r, R_\ell, R_b$ and $\Delta r \equiv 1 - \pi \alpha_{\text{EM}}(\sqrt{2} G_F)^{-1} M_W^2 (1 - M_W^2 M_Z^2)^{-1}$ from experiment. The data (at 90% c.l.) allow a band, as plotted against an MSSM mass parameter, say the mass of the CP-odd Higgs $m_A$ (Fig. 3). Three typical PMSSM fits are plotted \[15\]. The bold line corresponds to $\tan \beta = 70$, the dashed one to $\tan \beta = 20$, the dashed-dotted to $\tan \beta = 8$, the long-dotted to $\tan \beta = 1.5$ and the dotted one to $\tan \beta = 0.7$. The average squark and slepton masses have been fixed at 900 GeV and 500 GeV respectively while the gluino mass has been kept at 800 GeV. The $\mu$-parameter and $M_{1/2}$ have been kept at $-100$ GeV and 300 GeV respectively and no $L-R$ mixing has been assumed. Evidently, all quantities except $R_b$ and $R_c$ can be fit naturally and easily; however, in these MSSM does not necessarily do better than SM.

Focusing on $R_b$ and $R_c$, the experimental numbers quoted in conferences during the summer of ’95 are:

$R_b = 0.2219 \pm 0.0017$ (c.f. SM value 0.2157, i.e. $+3.7\sigma$ deviation),

$R_c = 0.1543 \pm 0.0074$ (c.f. SM value 0.171, i.e. $-2.5\sigma$ deviation).

Furthermore, assuming that $R_c$ is given by the SM value, one finds the experimentally determined $R_b$ to be $0.2206 \pm 0.0016$. No set of parameters within MSSM can fit $R_c$ consistently with other measured quantities. Ignoring $R_c$, several attempts have been made to fit $R_b$ through 1-loop enhancements. Broadly, there are two alternative scenarios: (1) the low $\tan \beta$, low mass ($< 90$ GeV) chargino and stop option or (2) the large $\tan \beta$, low mass CP-odd scalar $A$ option.

For (1) \[16\], one makes use of the one loop vertex correction with the supersymmetric analogue of the top-top-charged Higgs triangle, namely the stop-stop-chargino triangle. The couplings are proportional to $\cot^2 \beta$ and the low-mass stop and chargino enhance the propagator contributions to the second triangle. Thus $R_b$ can be enhanced. However, several cautionary remarks are in order. First, $\Gamma(b \to s \gamma)$ has to be kept under control – a task not easily done. Second, with such light stops, the decay $t \to \tilde{t} + \text{LSP}$ may become too large, creating problems with the CDF/D0 data. Finally, the lack of observation of any chargino at LEP 1.5 has poured cold water on this option. Claims have also been made that the $\alpha_S(M_Z)$ problem (namely, the discrepancy between its values determined from LEP and from lower energy measurements) is alleviated in this scenario, the SM $Z$-lineshape determination of $\alpha_S = 0.126 \pm 0.006$ being pulled down to $0.112 \pm 0.005$.

Turning to (2), the saviour \[17\] is the triangle with $b$, $b$ and $A$ internal lines. This will enhance the $A$-induced 1-loop contribution to $\Gamma(b \to c \tau \nu_\tau)$ which needs to be controlled. Moreover, one will have the decay $Z \to b\bar{b}A$ with the $A$ decaying further into $b\bar{b}$. This has not been seen. Using these two constraints, Wells and Kane \[17\] claimed to have excluded this option. However, there is the question of experimental efficiency in detecting a four $b$ final state in $Z$-decay. Taking that into account, the option is still viable \[18\].
Figure 3: Electroweak observables plotted against $m_A$
• **LEP 2 futurology**

The available CM energy of this machine will be in the $161 - 200$ GeV range. The dominant SM process will be $WW$ production for which the cross section has been calculated to be nearly $18$ pb, in comparison with a chargino pair production cross section (for a chargino that is nearly mass-degenerate with the $W$) in the vicinity of $5$ pb. This reaction will easily yield three types of final state configurations which are usually associated with sparticles: (1) acollinear lepton-pair with $E_T$, (2) acollinear jets with $E_T$ and (3) acollinear jets + lepton with $E_T$. Thus, $W$-pair production will provide a severe background to supersymmetry search at LEP 2. Nevertheless, there are three redeeming features \[14\]. First, the rates of the actual modes that are to be detected need to be calculated by folding branching ratios with the pair production cross section; after that is done, the non-supersymmetric channels do not fare particularly better over their supersymmetric rivals. Second the geometries of the two types of events differ significantly; those generated from $W$-pairs would tend to be anisotropic with a forward peak in contrast with the supersymmetric ones which are more isotropic. Finally, the kinematics are very different, the $W$ will largely have a 2-body decay with $E_T$ being associated with a massless particle whereas the chargino will have a 3-body decay with the $E_T$ coming from massive LSP; intelligently designed cuts will help discriminate between signal and background.

Another helpful procedure \[14\] will be a stage-by-stage increase in energy. This will keep the requisite number of open channels in check and will facilitate the search process. Thus, as $\sqrt{s}$ is increased, one expects first the pair-production of charginos $\tilde{\chi}^\pm$ and neutralinos $\tilde{\chi}^0$. These cross sections are expected to be in the several picobaran range. At a higher energy, one may expect right-chiral slepton pair production as the next process. Here the cross section per channel is anticipated to be somewhat smaller and therefore one needs to be able to distinguish the difference between that and the previous process. In fact, a comprehensive calculation of selectron pair-production at LEP 2 has been done \[20\] with care taken in delineating this boundary.

• **Bottomline**

In a nutshell, there is no direct evidence for weak-scale supersymmetry yet. Though the $R_{b,c}$ data are interesting, too much need not be read into them. In particular, any negative $\delta R_c$ has to go away first. LEP 2 will be a crucial probe on charginos. If no chargino is seen there, one might forego the “round the corner” option. Then a pessimist believer might expect supersymmetry to manifest itself through the “decoupling” scenario and just wait for the LHC.

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