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NEW PARTICLES AND INTERACTIONS*

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Introduction

Each one of us probably has different reasons for attending this meeting. My own interest in multiparticle dynamics is no longer as an end in itself, but as a means for advancing to the next stage in physics. I take the point of view (developed in Section 1) that the Standard SU(3) $\times$ SU(2) $\times$ U(1) Model is well established, and is now ripe to be used as a tool for analyzing physics beyond the Standard Model. The tool kit should include reliable and efficient ways of distinguishing gluon jets from quark jets, and of discriminating between t, b, c and light quark jets. In Section 2 I review what I consider to be the most topical physics issues arising from the recent confirmation of the Standard Model. These include the need for dynamical principles which go beyond the gauge principle, and in particular a satisfactory mechanism for gauge symmetry breaking. Some of the ideas proposed for solving these problems, such as technicolor and supersymmetry (SUSY), are reviewed in Section 3, together with some of the experimental tests that can be performed. Section 4 concentrates more deeply on SUSY and discusses some ways of looking for sparticles in e^+e^- and $\bar{p}p$ collisions.

In each of the Sections 2, 3 and 4, I try to emphasize the crucial role to be played by the multiparticle jet tools of Section 1 in resolving some hot physics issues. We will see in particular that the ability to discriminate heavy quark jets with high efficiency will be important, as will be good calorimetry and the ability to select (veto) events with (out) leptons. Finally Section 5 poses a question.

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1. Future Tools of the Trade

With the recent discoveries of the $W^{\pm}$[1] and of the Z^0,[2] all aspects of the Standard SU(3) $\times$ SU(2) $\times$ U(1) model can be consigned to the text books with the sole exception of the Higgs sector. We know that perturbation theory works well for the SU(2) $\times$ U(1) interactions – QED has been verified[3] through six orders of perturbation theory, and the successful predictions[4] of the $W^{\pm}$ and Z^0 masses include mass shifts of order 3 GeV due to one-loop weak radiative corrections. Unfortunately, as we have seen repeatedly at this conference, perturbation theory does not work so well as for the SU(3) strong interaction theory called QCD. It should be emphasized that QCD is a theory and not a model, and its problems are technical rather than fundamental. They are all related to the largeness of the strong coupling α_s, which is in turn related to our present closeness (on a logarithmic scale) to the Landau[5] pole Λ_3, the energy scale Q where the strong interactions appear in perturbation theory to blow up: $\alpha_s \propto 1/\ell n(Q^2/\Lambda_3^2)$ (see Fig. 1).

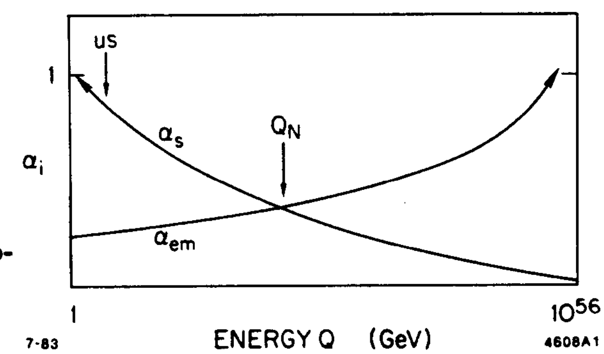


Fig. 1. Sketch of the logarithmic evolutions of the strong and electromagnetic couplings $\alpha_{s,em}$ illustrating their respective Landau poles, and a possible intermediate energy Nirvana Q_N where perturbation theory works well simultaneously for both interactions.

Life in QED is easy because we live a long way way from the electromagnetic Landau pole: $\alpha_{em} \sim 1/\ell n(\Lambda_{em}^2/Q^2)$ where $\Lambda_{em} = 0(10^{56})$ GeV. At these energies, on the other hand, the QCD coupling α_s would be so small that QCD perturbation theory would be very easy while QED peturbation theory would be very difficult. As indicated in Fig. 1, there may be an intermediate energy Nirvana Q_N: $\Lambda_3 \ll Q_N \ll \Lambda_{em}$ where both QED and QCD are easy to calculate. If so, we have not yet attained it! Despite these practical problems,[6] no serious theorist questions the status of QCD as the theory of the strong interactions. Thus QCD is not a fashion, though ways of calculating it which involve various model assumptions may be (are?) subject to changes in fashion.

The Standard SU(3) $\times$ SU(2) $\times$ U(1) Model contains the following established particles: quarks q, leptons ℓ, the photon γ, gluons g and intermediate vector bosons $W^{\pm}$ and Z^0. These are the building blocks which will populate the decays of particles yet to be discovered. Among them, the photon and the leptons are easy to identify experimentally, with the caveat that the τ is more difficult to spot than the e and the μ, and has indeed only been identified conclusively in e^+e^- collisions. We have seen recently[1,2] that the $W^\pm$ and Z^0 can be identified with some efficiency (3 to 8%) from their leptonic decay modes, but no one has yet picked out a clear $W^\pm$ or Z^0 signal in $\bar{q}q$ decays. Since quarks and gluons are confined we do not see them directly, but only in jets at high energies, which brings us to the topic of multiparticle dynamics.

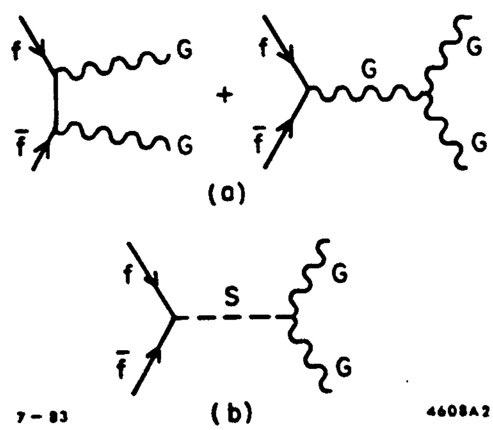
In an ideal world we would be able to identify and distinguish all the different flavors of q and g jets. In practice, all such distinctions can only be statistical, and no one yet knows how to draw a clear distinction even between q and g jets. The hope[7] that one might be able to separate u and d jets on the basis of their electromagnetic charges has long since been abandoned.[8] Also, it seems to be impossible to pick out strange quark jets by using the primary $s \to K$ or Λ fragmentation, since strange particles are also copiously pair-produced in u and d jets. Perhaps we should abandon attempts to distinguish among light quark jets. Several handles for picking out c quark jets are now available, though none of them is very efficient yet. High precision vertex detectors[9] can pick out charmed hadron decay paths. Many c jets yield a D^* which decays into $D\pi$ with a characteristic kinematic signature.[10] High resolution spectrometers can pick out a $D \to K\pi$ or $K\pi\pi$ bump, and c jets can be distinguished from lighter quark jets on the basis of semileptonic decays. Unfortunately, none of these identification tricks has more than a few percent efficiency at the present time, and the physics of the future may require a much higher tagging efficiency, as we will see shortly. Similar tricks may be applied to the identification of b quark jets. For example, it now seems possible that b-flavored hadrons may live long enough to be noticeable in a vertex detector,[11] and it has already proved possible[12] to pick out a b signal from jets containing leptons at large p_T. However, so far the efficiency for tagging b jets is also at the percent level. At one time it was hoped that one might be able to pick out b jets on the basis of their masses and/or their $\sum_{hadrons\ h} |p_T^h|$. Unfortunately, this hope has turned out to be illusory, but there are high hopes that it might be possible in the future to identify t jets by their "topology." Since $m_t \geq 20$ GeV, t jets should be quite wide and may break up into three identifiable subjets emanating from $t \to qq\bar{q}$ decay. The CERN $\bar{p}p$ collider may[13] soon be able to tell us whether these hopes for identifying t jets with high efficiency can be realized.

Let us assume that these jet identification tools can be honed to high efficiency, and go on to see how they may be used in resolving future physics issues.

2. Topical Physics Issues

Gauge invariance is now as sacred as motherhood and apple pie. Unfortunately, like these other precepts, it does not provide a complete framework for life. If weak SU(2) $\times$ U(1) gauge invariance were exact, it would enforce $m_q = m_\ell = m_W = m_Z = 0$. The non-zero values of these masses tell us that gauge invariance must be broken somehow. If we regard the retention of a valid perturbation theory for the weak interactions as a necessary feature of our physical theory, then theories with massive gauge bosons must contain spin-zero particles.[14] Perturbative unitarity requires that Figs. 2a must be supplemented by the scalar exchange of Figs. 2b, and that Figs. 3a must be supplemented by the scalar exchanges of Figs. 3b. Indeed, a complete analysis of perturbative unitarity to include external scalar particles reveals[14] that all the particle masses and couplings must be just those of a spontaneously broken gauge theory with Higgs fields. If perturbation theory is allowed to break down in the Higgs sector, then one is free to postulate dynamical symmetry breaking with composite Higgs bosons, as is done in the technicolor theories[15] to be mentioned in Section 3. If perturbation theory is followed all the way, then the spin-zero Higgs bosons must be elementary. Their couplings to spin 1/2 and spin 1 particles are related to their masses. The fermion masses are themselves without severe group-theoretical constraints and hence (almost) arbitrary. Ultimately there will be a need for some new principle beyond gauge invariance to constrain the Higgs-fermion couplings and fermion masses.

Fig. 2. When fermions (f) and gauge bosons (G) are massive, perturbative unitarity requires that the diagrams (a) must be supplemented by scalar (S) exchanges (b).



It is impossible to avoid either elementary or composite scalar fields, and in the minimal Standard Model there is[16,17] a single physical neutral Higgs boson H^0 with couplings that are completely specified:

$$\text{Fig. 2b} \quad \Rightarrow g_{H\bar{f}f} = \left(\sqrt{2}\,G_F\right)^{1/2} m_f \tag{1a}$$

$$\text{Fig. 3b}: \quad g_{HW^+W^-} = 2\left(\sqrt{2}\,G_F\right)^{1/2} m_W^2 \tag{1b}$$

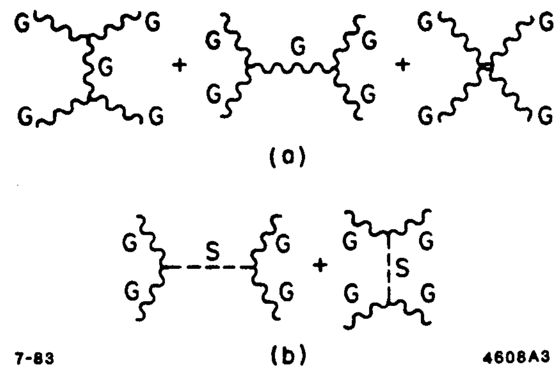


Fig. 3. Perturbative unitarity for massive gauge boson scattering requires that diagrams (a) be supplemented by scalar exchanges (b).

While the minimal version of the Standard Model contains no charged Higgs bosons $H^{\pm}$, they are present in all more complicated versions, and generally have couplings related to fermion masses in a manner less constrained than in Eq. (1a). Since all the other particles in the Standard Model have been found[13] in the minimal Standard Model, it only remains to hunt the Higgs.

<u>In e^+e^- Collisions</u> The branching ratio for toponium $\Theta \to H^0 + \gamma$ may be large[18]:

$$\frac{\Gamma(\Theta \to H^0 + \gamma)}{\Gamma(\Theta \to \gamma^* \to e^+e^-)} = \frac{G_F m_\Theta^2}{4\sqrt{2}\,\pi\alpha}\left[1 - \frac{m_H^2}{m_\Theta^2}\right] \; , \tag{2}$$

thanks to the large Higgs-$t\bar{t}$ coupling (1a). If no more dramatic large new decay modes for Θ open up, Eq. (2) suggests a branching ratio $\gtrsim 1\%$ for $m_\Theta \gtrsim 40$ GeV. The signature is a monochromatic photon accompanied by $H^0 \to \bar{b}b : \bar{c}c : \bar{\tau}\tau \simeq 10 : 1 : 1/3$ thanks to the mass-dependent couplings (1a). It would be important for pinning down the nature of the H^0 to be able to identify positively these heavy fermion decay modes, even with a sample of a few events. This may even be essential if one is to pick the $\Theta \to \gamma + H^0$ signal out from the $\Theta \to \gamma + gg$ background. Another favored reaction[19] for hunting the H^0 is $Z^0 \to H^0 + \ell^+\ell^-$:

$$\frac{\Gamma(Z^0 \to H^0 \to \ell^+\ell^-)}{\Gamma(Z^0 \to \ell^+\ell^-)} = 0(10^{-2\frac{1}{2}} \text{ to } 10^{-4\frac{1}{2}}) \; , \tag{3}$$

which gives a branching ratio $\approx 10^{-4}$ to 10^{-6} for $m_H = 10$ to 50 GeV. There are not going to be many $Z^0 \to H^0 + \ell^+\ell^-$ events, even with a Z^0 factory such as the SLC or LEP, and the ability to identify $H^0 \to \bar{b}\,b$, $\bar{c}\,c$ or $\bar{\tau}\,\tau$ decay modes with high efficiency will again be important. If $m_{H^0} > 0(50)$ GeV, the best way to find it may[20] be in the reaction $e^+e^- \to Z^0 + H^0$:

$$\frac{\sigma(e^+e^- \to Z^0 + H^0)}{\sigma(e^+e^- \to \gamma^* \to \mu^+\mu^-)} \geq 0.1 \quad , \tag{4}$$

for $m_H \leq 100$ GeV at a center-of-mass energy $\sqrt{s} \approx 200$ GeV. Such a heavy H^0 may have $\bar{t}\,t$ as its dominant decay mode. In view of the small cross-section (4) it may not be sufficient to use the $Z^0 \to \ell^+\ell^-$ decays with their small branching ratios, but one may need to use $Z^0 \to \bar{q}\,q$ and pick the reaction out from a multijet QCD background. We will have to understand and model QCD 3- and 4-jet reactions very thoroughly, and learn how to bring out the expected $H^0 \to$ heavy fermion decay signature.

The best way to look for the $H^\pm$ may be via $e^+e^- \to H^+H^-$:

$$\frac{\sigma(e^+e^- \to \gamma^* \to H^+H^-)}{\sigma(e^+e^- \to \gamma^* \to \mu^+\mu^-)} = \frac{1}{4}\beta^3 \tag{5}$$

and one might expect the $H^\pm$ to decay into two-jet $t\,\bar{b}$, $b\,\bar{c}$ or $c\,\bar{s}$ systems. However, these decay modes are more model-dependent (and hence informative and interesting) than the single H^0 decays discussed earlier. Again we will need to understand and fight down the QCD 4-jet background. Some tricks using jet angles, energies and invariant masses have already been developed by the TASSO Collaboration[21] in their pioneering search for $H^\pm \to$ hadronic jets.

<u>In *hh* Collisions</u> The basic reaction would be $hh \to H^0 + X$ which is expected to be dominated by gluon fusion.[22] The production cross section depends sensitively on the ratio between the t and H^0 masses. The best prospects for H^0 detection seem to be either when $m_{H^0} \approx 10$ GeV, in which case the rare decay mode $H^0 \to \mu^+\mu^-$ may peek out above the Drell-Yan background in a very high resolution experiment such as that planned[23] for the Tevatron, or else if $m_H > 200$ GeV in which case the dominant decay modes may[20] be $H^0 \to W^+W^-$ and Z^0Z^0 via the couplings (1b). In this latter case one probably needs to be able to pick out the $W^\pm$ and Z^0 with higher efficiency than is possible through their leptonic decay modes, if one is to have an observable event rate. Another suggested[24] reaction, suitable for $m_{H^0} < 0(50)$ GeV, is $hh \to W^\pm + X$ followed by $W^\pm \to H^0 + (\ell^\pm\nu)$. This decay has a similar branching ratio to that for $Z^0 \to H^0 +$

$\ell^+\ell^-$, but presumably it is more difficult to find because of the larger number of background hadrons in hh collisions. The kinematics force[25] the spectator final state $(\ell^\pm\nu)$ pair to have an invariant mass close to $m_W - m_H$, and hence $p_T^{\ell\pm} \approx (m_W - m_H)/2$. Some optimists talk about using very high luminosity hh colliders by throwing away all particles with $p_T < 0(5)$ GeV. In this way they would lose the H^0 decay products if m_{H^0} is small, and the spectator $\ell^\pm$ if m_{H^0} is large. A related reaction is $hh \to W^\pm$ (or Z^0)+$H^0 + X$, which has[26] a cross section $\leq 0(10^{-3})$ $\sigma(hh \to W^\pm$ (or Z^0) $+X)$ for $m_{H^0} > 10$ GeV. In this case one may benefit from being able to impose the kinematical constraint that $m_{\ell^\pm\nu}$ or $m_{\bar{q}q} = m_{W^\pm}$ ($m_{\ell^+\ell^-}$ or $m_{\bar{q}q} = m_{Z^0}$), but otherwise many of the same background worries apply to this reaction as to the previous one. The reaction $hh \to H^+H^- + X$ has a negligible cross section unless $m_{H^\pm} \leq m_{Z^0/2}$, and it is difficult to see how it could be detected against the overwhelming QCD jet background. This is another example of the problems in seeing non-strongly interacting particles in hh collisions. QCD jet backgrounds are relatively larger than in e^+e^- collisions by a factor $0(\alpha_3/\alpha)^2 \geq 0(10^2)$. Has anyone ever seen a τ in a hh collision?

My belief is that e^+e^- collisions are the best places to look for Higgs bosons (or whatever replaces them – see Section 3) and this is a large element in my enthusiasm for high energy e^+e^- colliders such as the SLC, LEP and beyond.[27]

3. New Ideas to be Tested

The standard Model with its elementary Higgs fields may appear satisfactory at first sight, but it has problems. In this section we first review these problems and then discuss possible solutions.

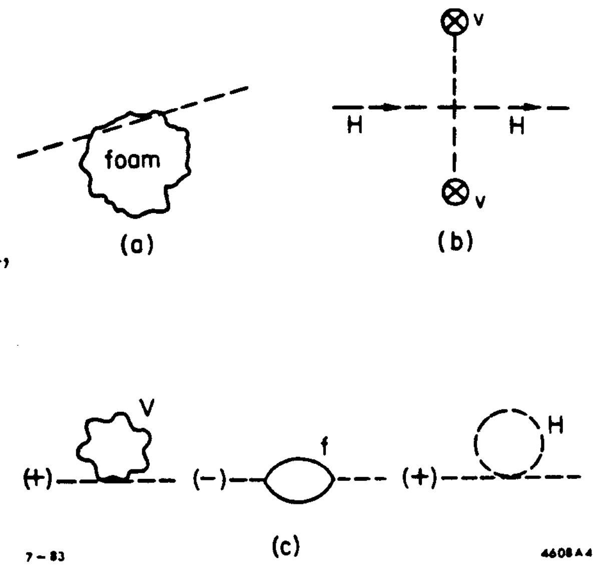
The Trouble With Higgs   We saw in Section 2 that if perturbation theory is to work for the weak interactions there must be at least one physical Higgs boson with $m_H = 0(m_W)$, and in any case $\leq 0(1)$ TeV. The trouble with this is that elementary scalar particle masses such as m_H are very unstable. Hawking and collaborators[28] have argued that elementary scalars propagating through the sort of foamy structure of space-time that one expects at the Planck length (Fig. 4a) will undergo mass shifts:

$$\delta m_H^2 = 0(m_P^2) \simeq (0(10^{19})\,\mathrm{GeV})^2 \quad . \tag{6}$$

Even if one can overcome this foamy hurdle, propagation through the grand unified theory (GUT) vacuum (Fig. 4b) (if you believe in GUTs) will give

$$\delta m_H^2 = 0(m_X^2) = (\geq 0(10^{15})\,\mathrm{GeV})^2 \tag{7}$$

Fig. 4. An elementary scalar particle can acquire a large mass (a) from propagating through space-time foam, (b) from propagating through the GUT vacuum, or (c) from radiative corrections.



from interactions between the heavy Higgses with $m_H = 0(m_X) \geq 0(10^{15})$ GeV and those that are trying to be light with $m_H = 0(m_W)$. Even if one could see how to make the large contributions (6, 7) to m_H^2 cancel or vanish (why?), radiative corrections such as those in Fig. 4c would undo the cancellations[29] and yield

$$\delta m_H^2 = 0(\alpha^n) \times 0(\Lambda^2) \ , \tag{8}$$

where Λ is a cut-off on the momenta circulating in the loops of Fig. 4c, which could be as large as $\Lambda = 0(m_X$ or $m_P)$. For m_H to be $0(m_W)$ we must find some way of eradicating these radiative corrections. This could be done by dissolving the perturbation theory diagrams at internal momenta $\Lambda = 0(1)$ TeV corresponding to a scale at which the spin-zero fields are composite, which is the approach followed in theories[15] of dynamical symmetry breaking (technocolor). In these theories symmetries prevent some composite spin-zero bosons from acquiring masses, while others acquire $\delta m^2 = 0(\alpha^n)0(\Lambda^2 = 0(1)\,\mathrm{TeV}^2)$. If the scalar fields are not composite on some such small scale Λ, we must cancel the diagrams of Fig. 4c through order α^{12} or α^{16} if there is no cut-off before $\Lambda \approx m_X$ or m_P. This can be done by adding fermion and boson diagrams which are considerate enough to have opposite signs. If there are bosons (+) and fermions (−) with identical couplings, one-loop diagrams give

$$\Delta m_H^2 = 0(\alpha)|m_B^2 - m_F^2| \ . \tag{9}$$

Identical couplings means supersymmetry (SUSY).[30] Getting Eq. (9) to be 0(100 GeV)^2 requires supersymmetric partners X and $\tilde{X}$ to have similar masses

$$|m_{\tilde{X}}^2 - m_X^2| \leq 0(1)\,\mathrm{TeV}^2 \tag{10}$$

so that SUSY must be a "good" approximate symmetry valid at scales $\ll m_X$ or m_P.

Let us now discuss these new ideas of technicolor and supersymmetry in greater detail.

Technicolor The central feature of this theory[15] is the existence of a new set of strong gauge interactions which confine technifermions on a distance scale $d = 0(1/\Lambda_T) = 0(1/1\,\mathrm{TeV})$ in much the same way as QCD confines quarks on a scale $d = 0(1/\Lambda_3) = 0(1/1\ \mathrm{GeV})$. The conventional elementary Higgs is replaced by a bound state technipion π_T analogous to the conventional pion π:

$$H \rightarrow \pi_T = (\bar{F}\,F)_T \leftrightarrow \pi = (\bar{q}\,q) \tag{11a}$$

while the normal Higgs vacuum expectation value is replaced by a technifermion condensate in the vacuum

$$\langle 0|H|0\rangle \rightarrow \langle 0|(\bar{F}_L F_R)_T|0\rangle = 0(\Lambda_T^3) \leftrightarrow \langle 0|\,\bar{q}_L\,q_R|\rangle = 0(\Lambda_3^3) \tag{11b}$$

which breaks the weak gauge symmetry spontaneously. The $W^\pm$ and Z^0 eat massless technipions as in Fig. 5 and thereby acquire masses whose scale is fixed dynamically:

$$m_{W^\pm} = 0(g_2)\Lambda_T \quad . \tag{12}$$

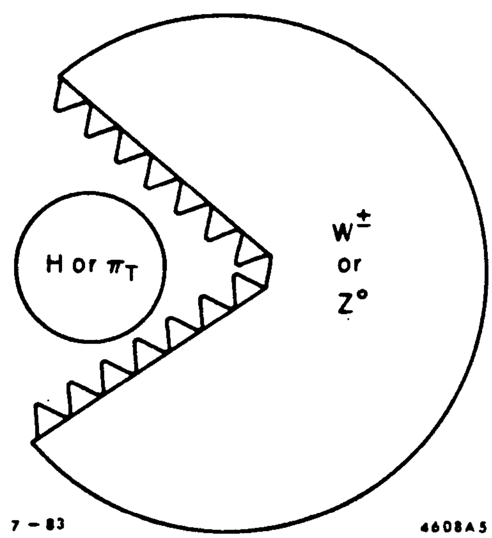


Fig. 5. The mechanism of mass generation for the $W^\pm$ and Z^0.

It should eventually be possible to calculate Λ_T when technicolor is unified[31] with the other interactions analogously to the way QCD is embedded into GUTs along with the weak and electromagnetic interactions

$$\Lambda_{QCD} = m_X \exp\left(\frac{-0(1)}{\alpha}\right) \rightarrow \Lambda_T = m_{X'} \exp\left(\frac{-0(1)'}{\alpha}\right) \quad . \tag{13}$$

This scenario generates $m_{W^\pm}$ and m_{Z^0} very economically, but it requires unaesthetic epicycles if one is to generate masses for the conventional quarks and leptons. One possibility[32] is to add in extended technicolor (ETC) interactions mediated by massive gauge bosons whose exchanges as in Fig. 6 yield

$$m_{q_1\ell} = 0\left(\frac{1}{m_{ETC}^2}\Lambda_T^3\right) \quad . \tag{14}$$

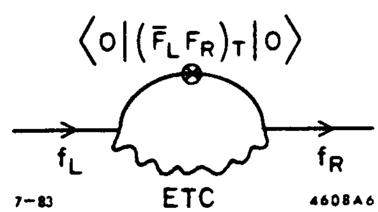


Fig. 6. The ETC mechanism[32] of mass generation for conventional fermions.

In addition to its inelegance (we see from the formula (14) that many different ETC bosons with different masses may be required for generating all the masses of the different quark and lepton generations) the ETC scenario has various experimental problems.[33] One is that the exchanges of heavy particles akin to the ETC bosons give rise to flavor-changing neutral interactions whose magnitudes are estimated[34] to exceed present experimental upper limits. A second problem is that the simplest ETC models predict[35] the existence of uneaten charged technipions $P^\pm$ with masses

$$m_{P^\pm} \approx 0(5 \text{ to } 14) \text{ GeV} \tag{15}$$

which are almost excluded by experiments[21,36] at PEP and PETRA. Although there is no complete technicolor model which avoids all these problems, it is not impossible that one might be constructed.[37] and therefore experimentalists should continue looking for technicolor.

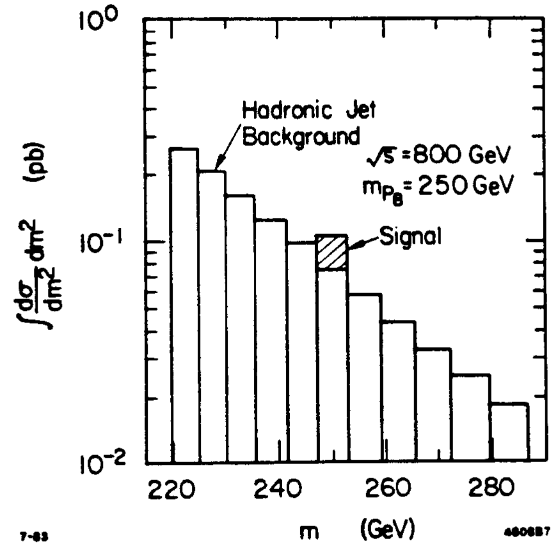
<u>In e^+e^- Collisions</u> The reaction $e^+e^- \rightarrow P^+P^-$ is very similar to the reaction $e^+e^- \rightarrow H^+H^-$ discussed in Section 2. Again one might expect $P^\pm \rightarrow t\,\bar{b}$, $b\,\bar{c}$ and $c\,\bar{s}$ decays to dominate, and specific models enable definite predictions to be made

for the different branching ratios.[38] These render the present absence of the $P^{\pm}$ particularly disturbing. They also highlight the interest of being able to identify efficiently heavy quark jets in the future: these should indeed be the dominant decay modes of spin-zero particles, and detailed measurements of the branching ratios may distinguish between different models. A difference between elementary Higgs theories and technicolor models is that composite neutral scalars P^0 are not expected[38,39] to have large couplings to the Z^0, and hence the branching ratio for $Z^0 \rightarrow P^0 + \ell^+\ell^-$ and the cross section for $e^+e^- \rightarrow Z^0 + P^0$ are both expected to be unobservably small. On the other hand toponium $\Theta \rightarrow P + \gamma$ is expected[38] to occur with a branching ratio within a factor three of that for $\Theta \rightarrow H^0 + \gamma$ (2). The new strong interactions of technicolor theories offer other exciting prospects for e^+e^- experiments at high energy. For example, there should be a techni-ρ which shows up as a significant bump[40] in the total e^+e^- annihilation cross section around 1 TeV, and should have a large branching ratio into W^+W^- pairs. There should also be a techni-η' with mass 0(1) TeV which could be seen in $\gamma\gamma$ collisions at a high energy e^+e^- collider.[27]

<u>In *hh* Collisions</u> Production of the P^0 and $P^{\pm}$ parallels that of the corresponding Higgses in some respects: $hh \rightarrow (gg \rightarrow P^0) + X$, $hh \rightarrow (\gamma^*$ or $Z^0 \rightarrow P^+P^-) + X$. However, their small couplings to the gauge bosons mean that production of the P^0 or $P^{\pm}$ in association with real or virtual vector bosons is negligible. There are other heavier technicolor particles whose production might be observable in high energy hh collisions. These include color octet pseudoscalars P_8: $hh \rightarrow (gg \rightarrow P_8) + X$, which are expected to have masses about 250 GeV and decay mainly into $\bar{t}\,t$ and gg jets. This is another reason to be able to identify these jets with high efficiency, particularly in view of the low signal-to-background ratio when a plausible dijet mass resolution is taken into account[41] as in Fig. 7. At even higher masses, there is a large number of natural J^{PC} technicolor resonances which can be produced in hh collisions and subsequently decay into pairs of technipions, notably including the longitudinal polarization states of the $W^{\pm}$ and Z^0. To pick them out, one may again require the efficient identification of vector bosons decaying into $\bar{q}\,q$ jets.

I am not convinced that the prospects for technicolor searches in hh collisions are much brighter than those for conventional Higgs searches.

Fig. 7. A calculation[41] of colored technipion P_8 production in hadron-hadron collisions, comparing its $P_8 \to \bar{t}\,t$ dijet decay signature with the QCD background, assuming a plausible dijet mass resolution.



Supersymmetry This is a new kind of symmetry which relates fermions to bosons.[30] It is generated by spin-1/2 charges Q_α which, as one might expect of fermionic operators, obey an anticommutation algebra:

$$\{Q^i_\alpha\ ,\ Q^{+\dot{\alpha}}_i\} = 2(\sigma^\mu)^{\dot{\alpha}}_\alpha P_\mu \quad . \tag{16}$$

In addition to the spinorial index α, the charges leave an internal symmetry index $i = 1, 2, \ldots, N$. The case $N = 1$ is known as simple supersymmetry, while theories with $N > 1$ are said to possess extended supersymmetry. Particles in gauge theories are restricted to helicities $\lambda : |\lambda| \leq 1$. Since each supersymmetry transformation Q changes the helicity by half a unit, one can tolerate at most $N = 4$ in a gauge theory:

$$\lambda = +1 \ \underset{Q}{\rightarrow}\ +\frac{1}{2}\ \underset{Q}{\rightarrow}\ 0\ \rightarrow\ -\frac{1}{2}\ \underset{Q}{\rightarrow}\ -1\ . \tag{17}$$

On the other hand, if one includes gravity: $|\lambda| \leq 2$ then one must make supersymmetry a local symmetry, and then can tolerate $N \leq 8$. We will restrict ourselves to simple supersymmetry: $N = 1$ in the phenomenological discussion that follows. In this case the only possible supermultiplets are

$$\begin{aligned} &\text{gauge}: && \lambda = \left(1, \frac{1}{2}\right) \ ; \\ &\text{chiral}: && \lambda = \left(\frac{1}{2}\ , 0\right) \end{aligned} \tag{18}$$

as well as the graviton supermultiplet $\lambda = (2,\ 3/2)$.

Unfortunately, no known particle can be the spartner of any other known particle. This means that we must double* the known set of helicity states, as seen in Table 1.

Table 1. Particles and Their Supersymmetric Partners

Particle	Helicity	Sparticle	Helicity
quark q	1/2	squark $\tilde{q}$	0
lepton ℓ	1/2	slepton $\tilde{\ell}$	0
photon γ	1	photino $\tilde{\gamma}$	1/2
gluon g	1	gluino $\tilde{g}$	1/2
vector bosons $W^{\pm}$	1	wino $\tilde{W}^{\pm}$	1/2
vector bosons Z^0	1	zino $\tilde{Z}^0$	1/2
Higgs H	0	shiggs $\tilde{H}$	1/2

Since the charged sparticles can be pair-produced copiously in e^+e^- annihilation, and none has yet been seen, they must all have masses

$$m_{\tilde{q}}\ ,\ m_{\tilde{\ell}}\ ,\ m_{\tilde{W}^{\pm}}\ ,\ m_{\tilde{H}^{\pm}} \geq 0(17)\ \text{GeV} \quad . \tag{19}$$

Neutral particles are more difficult to produce, but colored ones could already have been seen in hh collisions, and have not been, so that[43]

$$m_{\tilde{g}} \geq 0(2\ \text{to}\ 3)\ \text{GeV} \quad . \tag{20}$$

On the other hand, there is no particle physics reason why a colorless neutral sparticle should not have a negligibly small mass:

$$m_{\tilde{\gamma}}\ ,\ m_{\tilde{H}^0} \to 0 \quad . \tag{21}$$

The lightest sparticle is essentially stable in most models, in which case cosmology requires it to be neutral. This has important implications for the preferred experimental signatures of supersymmetry which are discussed in the next section.

4. Search for SUSY

The most favored possibility[44] is that the lightest sparticle is a mixture of neutral gauginos and higgses which is almost a pure photino $\tilde{\gamma}$. There are model-building reasons why one might expect gravitinos or pure shigges to be heavier,

*In fact, more than double, since the SUSY Standard Model has at least two Higgs doublets $H_{1,2}$.

and cosmology certainly prefers[44] $m_{\tilde{H}} > m_{\tilde{\gamma}}$ in a wide class of models. In this case the preferred experimental signature of sparticle production and decay is

$$\tilde{X} \to X + \tilde{\gamma} \;:\; X = q, \ell, W^{\pm}, \ldots \tag{22}$$

where the $\tilde{\gamma}$ escapes the detector as missing neutral energy analogously to a neutrino, with the exception that comology strongly suggests[44,45]

$$m_{\tilde{\gamma}} \geq 0\left(\frac{1}{2}\right) \text{ GeV} \quad . \tag{23}$$

An alternative possibility[44] is the existence of a very light ($m_{\tilde{H}} \leq 0(100)$ eV) neutral shiggs, which would usurp the role (22) of the $\tilde{\gamma}$ in sparticle decays. In either case similar experimental considerations apply.

There is a premium placed on calorimetry, so as to be able to detect that missing energy was carried off as in (22). This calorimetric capability should be combined with the ability to veto on the presence of leptons in the final state. This is to discriminate between heavy quark decay $q \to q'\ell\nu$ and squark decay $\tilde{q} \to q + \tilde{\gamma}$ (or $q + \tilde{H}$). Jetology is also necessary if one is to be able to discriminate between $\tilde{t}$, $\tilde{b}$, $\tilde{c}$, $\tilde{s}$, $\tilde{u}$ and $\tilde{d}$, all of which must have similar masses in many models.[46]

Sample searches[42] for sparticles include hunts for charged sleptons assuming the production and decay pattern:

$$\begin{array}{l} \qquad\qquad\quad \ulcorner\!\!\to \ell^- + \tilde{\gamma} \\ e^+e^- \to \tilde{\ell}^+\tilde{\ell}^- \\ \qquad\qquad\quad \llcorner\!\!\longrightarrow \ell + \tilde{\gamma} \end{array} \tag{24}$$

which yields dilepton final states that are highly acoplanar with a large amount of missing energy. The same pattern has been sought for squarks:

$$\begin{array}{l} \qquad\qquad\quad \ulcorner\!\!\to \bar{q} + \tilde{\gamma} \\ e^+e^- \to \tilde{q}\,\tilde{\bar{q}} \\ \qquad\qquad\quad \llcorner\!\!\to q + \tilde{\gamma} \end{array} \tag{25}$$

yielding acoplanar dijet events. However, the signature (25) may be diluted if the decay $\tilde{q} \to q + \tilde{g}$ is kinematically accessible in which case

$$\frac{\Gamma(\tilde{q} \to q + \tilde{g})}{\Gamma(\tilde{q} \to q + \tilde{\gamma})} = (3 \text{ or } 12) \left(\frac{\alpha_s}{\alpha_{em}}\right) \gg 1 \ . \tag{26}$$

The $\tilde{g}$ would probably then decay via $\tilde{g} \to q\,\bar{q} + \tilde{\gamma}$ in which case the events would not have a simple two-jet structure and would have less missing energy than in the case (25).

An interesting way[47] to search for photinos directly is via the radiative annihilation reaction $e^+e^- \to \gamma + (\tilde{\gamma}\,\tilde{\gamma})$. The cross section for this process has a form very similar to that[48] for $e^+e^- \to \gamma + (\bar{\nu}\,\nu)$:

$$\frac{d^2\sigma}{dx_\gamma d(\cos\theta_\gamma)} = \frac{E^2_{c.m.}}{x_\gamma \sin^2\theta_\gamma} \left[(1 - x_\gamma)\left(1 - \frac{x_\gamma}{2}\right)^2 + \frac{x^2_\gamma}{4}(1 - x_\gamma)\cos^2\theta_\gamma\right]$$

$$\times \left\{\frac{G^2_F\alpha}{6\pi^2}\left[N_\nu(g^2_V + g^2_A) + 2(g_V + g_A + 1)\right]\right\} \ : \ g_V = \frac{1}{2} = -g_A \ \text{for } \bar{\nu}\,\nu \tag{27}$$

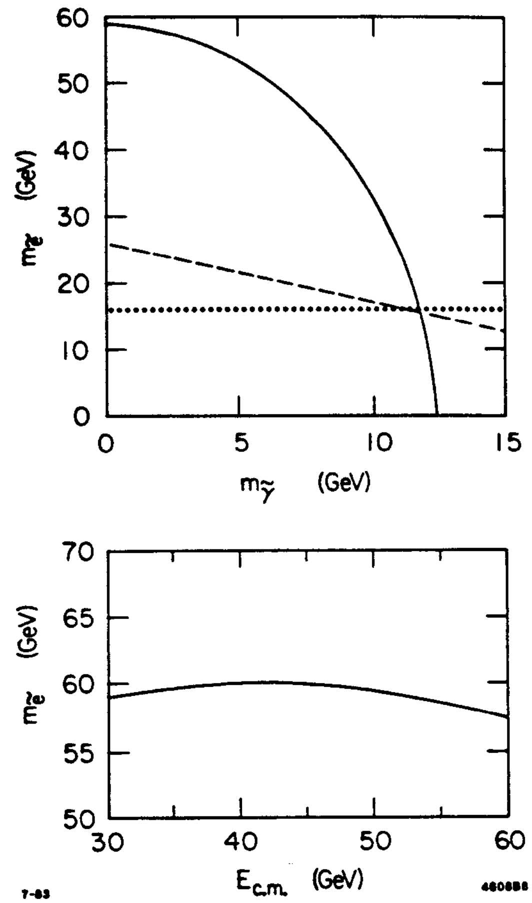
$$\text{or} \left\{\frac{2}{3}\frac{\alpha^2}{m^4_{\tilde{e}_{L,R}}}\right\} \text{ for } \tilde{\gamma}\,\tilde{\gamma} \ .$$

We have incorporated e^+e^- annihilation to $\tilde{\gamma}\,\tilde{\gamma}$ by selectron $\tilde{e}$ exchange without making any assumption about the ratio of masses of spartners of the left- and right-hand electrons $m_{\tilde{e}_{L,R}}$. Figure 8a demonstrates the sensitivity of this reaction to

$$m_{\tilde{e}} \leq 0(60) \text{ GeV} \ , \ m_{\tilde{\gamma}} \leq 0(10) \text{ GeV} \tag{28}$$

in an experiment with sensitivity to a cross section of order 10^{38} cm^2 at $E_{c.m.} =$ 29 GeV. Figure 8a also shows the ranges of $m_{\tilde{e}}$ and $m_{\tilde{\gamma}}$ that are accessible to $e^+e^- \to \tilde{e}^+\,\tilde{e}^-$ and $e^+e^- \to e^\pm\,\tilde{e}^\mp\,\tilde{\gamma}$ searches at present energies. Because of the background to $e^+e^- \to \gamma + (\tilde{\gamma}\,\tilde{\gamma})$ coming from $e^+e^- \to (\bar{\nu}\,\nu) + \gamma$, one does not gain much in sensitivity to $m_{\tilde{e}}$ as $E_{c.m.}$ is increased, as can be seen in Fig. 8b. To establish an upper limit on the cross section for $e^+e^- \to \gamma +$ nothing visible of order 10^{-38}cm^2 requires a detector which not also has good photon detection down to small energies and angles but also has no holes, so as to be able to reject backgrounds from such conventional QED sources as $e^+e^- \to \gamma =$ $(\gamma\gamma)$ or $\gamma + (e^+e^-)$. A recently approved experiment designed expressly with this intention is PEP-021.[49]

Fig. 8. Ranges of $m_{\tilde{e}}$ and $m_{\tilde{\gamma}}$ accessible[47] to present-day e^+e^- searches assuming $E_{c.m.} = 30$ GeV and $\sigma = 10^{-38}$ cm^2 for the reactions $e^+e^- \to \gamma + (\tilde{\gamma}\,\tilde{\gamma})$ and $e^+e^- \to e^\pm\, \tilde{e}^\mp\, \tilde{\gamma}$. (b) Change in sensitivity to $m_{\tilde{e}}$ as $E_{c.m.}$ is increased, assuming $m_{\tilde{\gamma}} = 0$.



There are other ways[50,51,52] to look for light gauginos, in particular at the CERN $\bar{p}\,p$ collider. In many models one or even two decay modes:

$$W^\pm \to \left(\chi^\pm \equiv \tilde{W}^\pm \text{ or } \tilde{H}^\pm\right) + (\chi^0 \equiv \tilde{\gamma} \text{ or } \tilde{H}) \tag{29}$$

are kinematically accessible,[53] and can have branching ratios comparable with the familiar $W^\pm \to e^\pm + \nu$ decay. Cosmology requires that the lightest charged gauge fermions $\chi^+ \equiv \tilde{W}^\pm$ or $\tilde{H}^\pm$ be heavier than the lightest neutral gauge fermion $\chi^0 \equiv \tilde{\gamma}$ or $\tilde{H}$. The charged gauge fermion will therefore be able to decay:

$$\chi^\pm \to \chi^0 + (\bar{q}\,q \text{ or } \bar{\ell}\,\ell) \ . \tag{30}$$

In general, the gauginos and shigges will mix. Their mass eigenstates can be obtained[50] by diagonalizing their mass matrices which are fixed by the parameters ϵ and M_2:

$$\mathcal{L} \ni \epsilon\ \epsilon_{\alpha\beta}\left(\tilde{H}_1^\alpha\, \tilde{H}_2^\beta\right) - M_2\, \tilde{W}_a\, \tilde{W}_a - M_1\, \tilde{B}^2 \tag{31}$$

(α, β are doublet, while a, b are triplet SU(2) indices) where

$$M_1 \approx \frac{5}{2} \frac{\alpha_1}{\alpha_2} M_2 \tag{32}$$

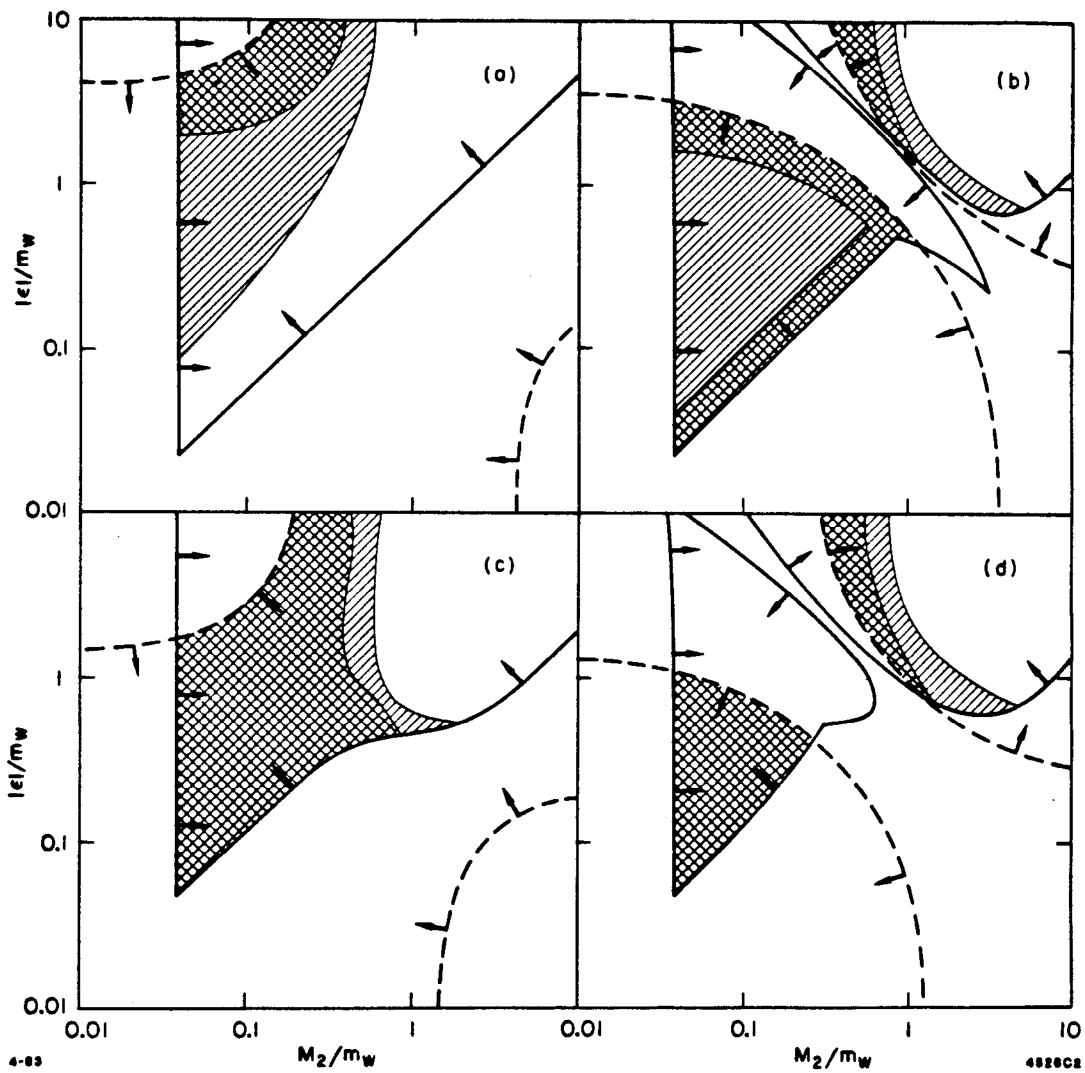


Fig. 9. Domains of ϵ, M_2 parameter space consistent[50] with cosmology (solid) and PEP/PETRA limits on $\chi^{\pm}$ (dashed). Also shown are the domains in which one (hatched) and two (cross-hatched) $W^{\pm} \rightarrow \chi^{\pm} + \chi^0$ decay modes are kinematically allowed for the following cases: (a) $v_1 = v_2$, $\epsilon > 0$; (b) $v_1 = v_2$, $\epsilon < 0$; (c) $v_1 = 4v_2$, $\epsilon > 0$; (d) $v_1 = 4v_2$, $\epsilon < 0$.

in models where SU(2) $\times$ U(1) is eventually unified in a GUT group. The parameters ϵ and M_2 are generally expected to be $0(m_W)$. Figure 9 shows the cosmologically allowed domains of ϵ and M_2 for a set of models characterized by different ratios of $v_{1,2} \equiv \langle 0|H_{1,2}|0\rangle$ and signs of the mixing parameter ϵ (31). Also shown is the constraint that the lightest charged gauge fermion have a mass of at least 20 GeV. The supersymmetric $W^{\pm}$ decays (29) are kinematically accessible in the shaded regions.

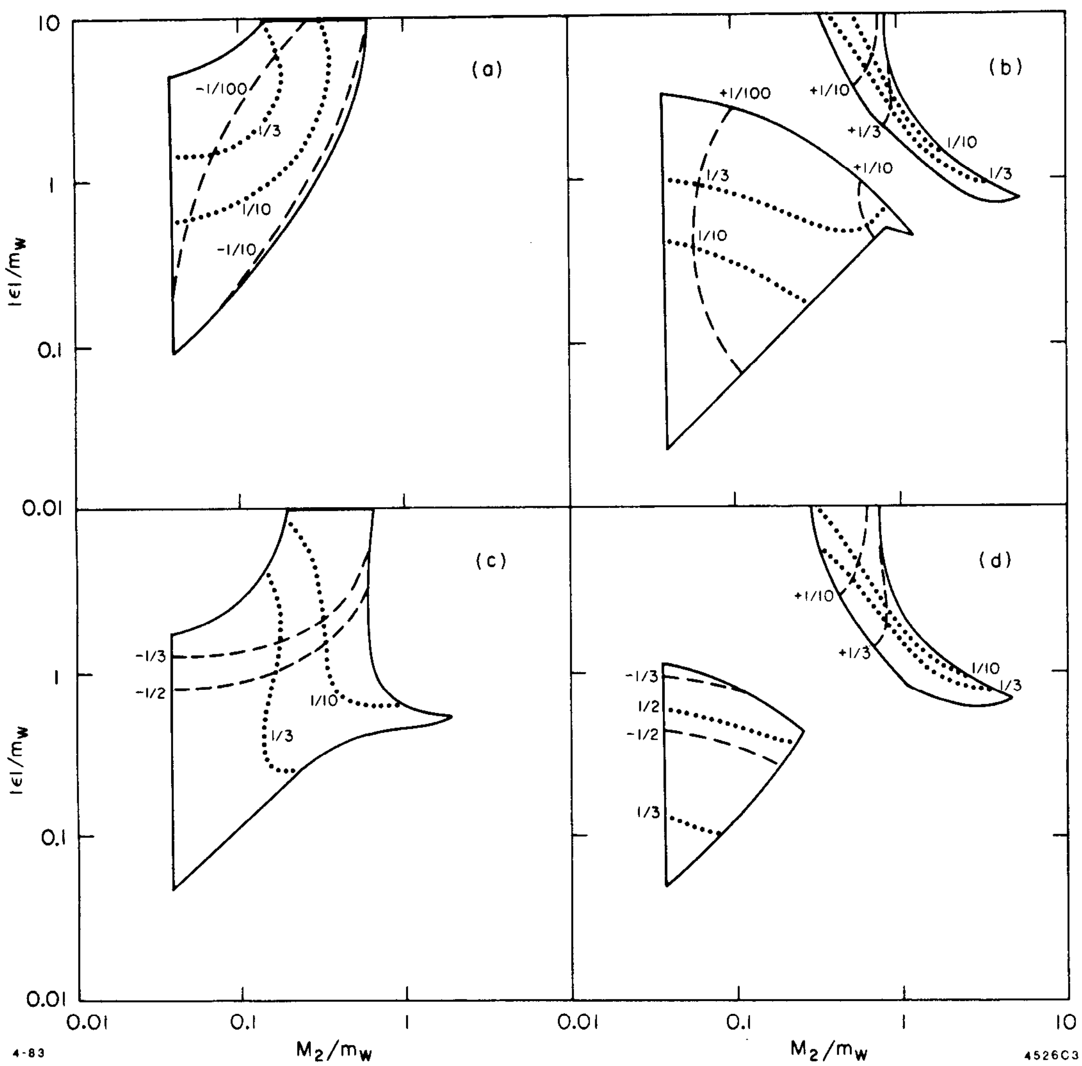


Fig. 10. Rates and forward-backward asymmetries[50] for "zen" events in the allowed regions of Fig. 9. Dotted lines are rates normalized to the $e\nu$ rate. Dashed lines represent forward-backward asymmetries. The labels (a) to (d) correspond to those in Fig. 9.

Figure 10 shows the rates for these decays as a fraction of the familiar $W^{\pm} \to e^{\pm}\nu$ decay rate, as well as the forward-backward decay asymmetry. Possible experimental signatures for the decays (29) are shown in Fig. 11. If the gauge fermions are light, the visible decay products (30) may be collimated so as to form a single jet system on one side of the collision axis: the "zen" signature[50] seen in Fig. 11a. If the charged gauge fermion is heavier, $\bar{q}\,q$ jets from its decay (30) will be more splayed out as in Fig. 11b, giving two jet events with a noticeable p_T imbalance.

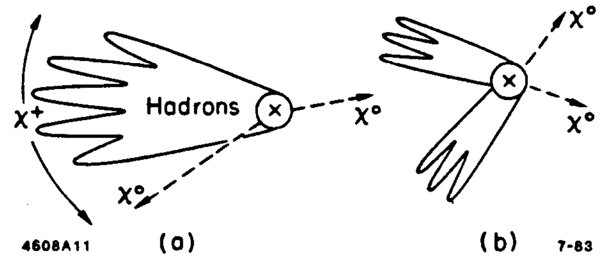


Fig. 11. Possible event signatures for sparticle production in $\bar{p}\,p$ collisions: (a) single jet "zen" events, and (b) two-jet events which are not back-to-back.

"Zen" events can also be produced in e^+e^- annihilation,[54,55] thanks to associated production $e^+e^- \to \chi^0\chi^{0'}$, where the χ^0 and the $\chi^{0'}$ are the lightest and the second lightest neutral gauge fermion respectively, and the $\chi^{0'}$ can decay:

$$\chi^{0'} \to \chi^0 + (\bar{q}\,q \text{ or } \ell^+\ell^-) \ . \tag{33}$$

In many models the $(\ell^+\ell^-)$ decay mode is favored.[55] The cross-section for $e^+e^- \to \chi^0\chi^{0'}$ may be comparable to conventional $e^+e^- \to \bar{\nu}\,\nu$ annihilation for certain values of the parameters ϵ, M_2 and light $\tilde{e}$ masses.[55] The pair-production reaction $e^+e^- \to \chi^{0'}\chi^{0'}$ can also have a cross-section[55] comparable to that for $e^+e^- \to \bar{\nu}\,\nu$ and would give 4-jet (or 2-jet and 2-lepton, or 4-lepton) events with missing energy.

Supersymmetric theories offer many other possible sources of "zen" events in e^+e^- annihilation and elsewhere. For example $\bar{p}\,p$ or $e^+e^- \to \tilde{\nu}\,\tilde{\nu}$ followed by a visible decay of just one of the sneutrinos $\tilde{\nu}$ may also have a rate comparable to that for $\bar{p}\,p \to Z^0$ or $e^+e^- \to \bar{\nu}\,\nu$.[51] The event signatures of Fig. 11 could also arise from the pair-production of other varieties of supersymmetric particles in $\bar{p}\,p$ collisions. Calculations of $hh \to (\tilde{q}\,\tilde{q},\ \tilde{q}\,\tilde{g}$ or $\tilde{g}\,\tilde{g}) + X$ all give cross-sections (see Fig. 12) which should be observable at the CERN $\bar{p}\,p$ collider if $m_{\tilde{q}},\ \tilde{g} \lesssim 0(50)$ GeV. In all these cases, missing transverse energy is a handy event signature. Calculations[57] of the signal and background for $\tilde{g}\,\tilde{g}$ production are shown in Fig. 13: the distributions in both the transverse momentum balance variable

$$x_E \equiv \frac{-\not{p}_{Jet_1} \cdot \not{p}_{Jet_2}}{|\not{p}_{Jet_1}|^2} \tag{34a}$$

and in the variable

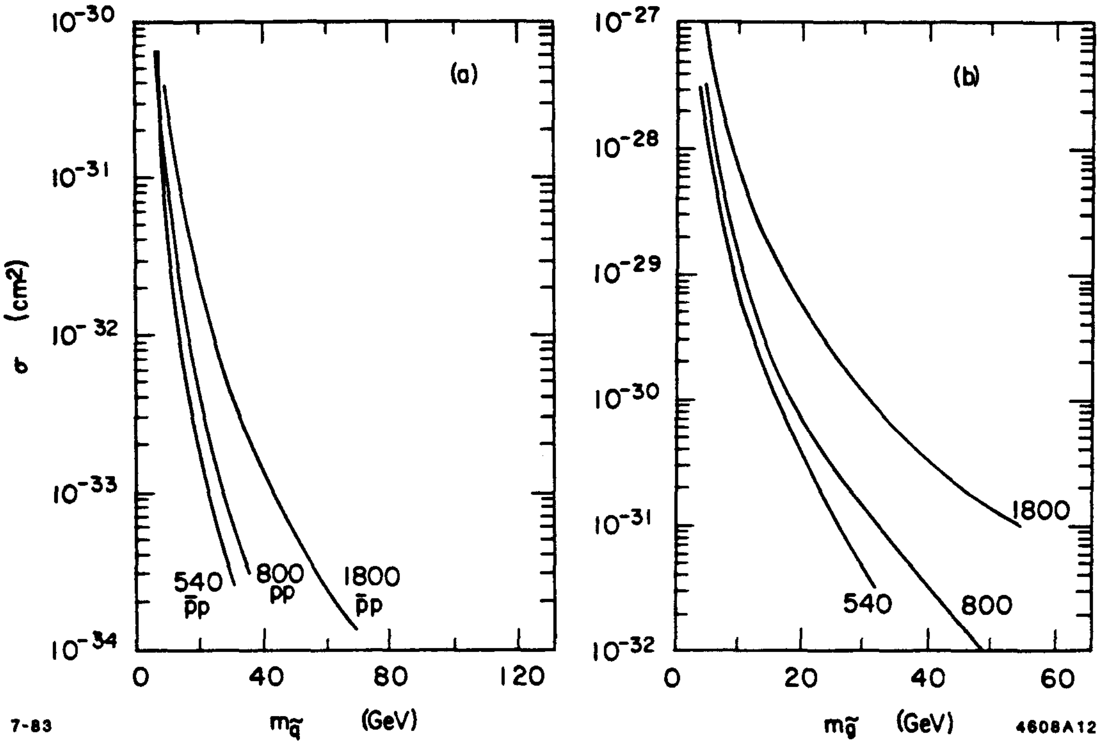


Fig. 12. Cross sections (a) for $\tilde{q}\,\tilde{q}$, and (b) for $\tilde{g}\,\tilde{g}$ production in hh collisions, taken from Kane and Leveillé, Ref. 43.

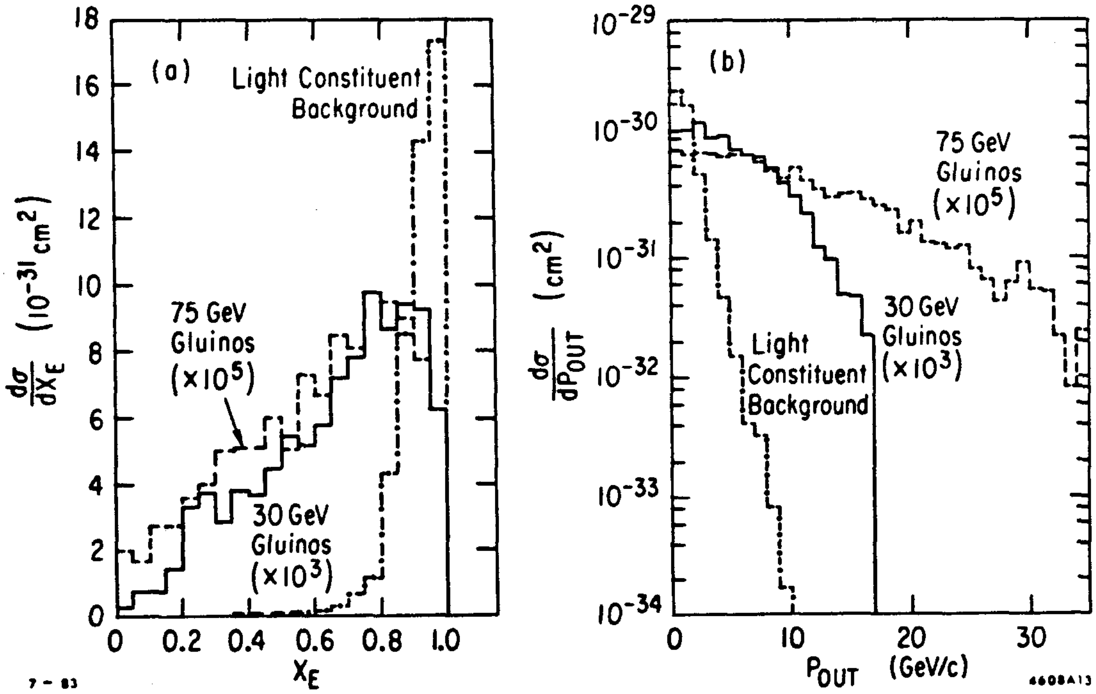


Fig. 13. Calculations[57] of the signals and backgrounds to sparticle production in hh collisions (a) in x_E, and (b) in p_{out}.

$$p_{out} \equiv \sqrt{|p_{Jet_2}|^2 - x_E^2 |p_{Jet_1}|^2} \tag{34b}$$

are shown.

It appears from these studies that there are several promising avenues for SUSY searches at present accelerators which have not yet been fully explored. Many of these studies require an understanding of jets or of other aspects of multiparticle dynamics, either to beat down the background or else pick out and analyze the signal. Figure 14 shows a short message from us theorists which may help our experimental colleagues to get motivated to search for SUSY. As reinforcement, Table 2 is a little historical reminder. Perhaps multiparticle dynamics will help usher in the SUSY revolution?

Table 2. Past and Future History

Gauge Theories		Supersymmetry	
Date	Event	Date	Superevent
1954	Invention of gauge theory (Yang, Mills)	1973	First SUSY theory (Wess, Zumino)
1961±	Early models (Glashow, etc.)	1977±	Early models (Fayet etc.)
1967-8	Standard Model (Weinberg, Salam)	?	?
1971	Renormalizability proven ('t Hooft)	1981	Application to hierarchy problem
1972	Searches for neutral currents	1983	This talk
1973	Discovery of neutral currents	?	?
1974	Discovery of charm	?	?
	etc.		etc.
the gauge revolution		the SUSY revolution	

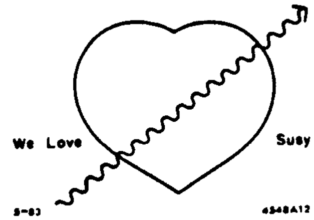


Fig. 14. A short message from theorists.

5. A Question

Figure 15 is an impression of how elementary particle physics is developing. Our subject searches for new laws and so addresses more fundamental questions as time progresses, whereas many other disciplines make more detailed studies of systems whose basic physical laws are understood.

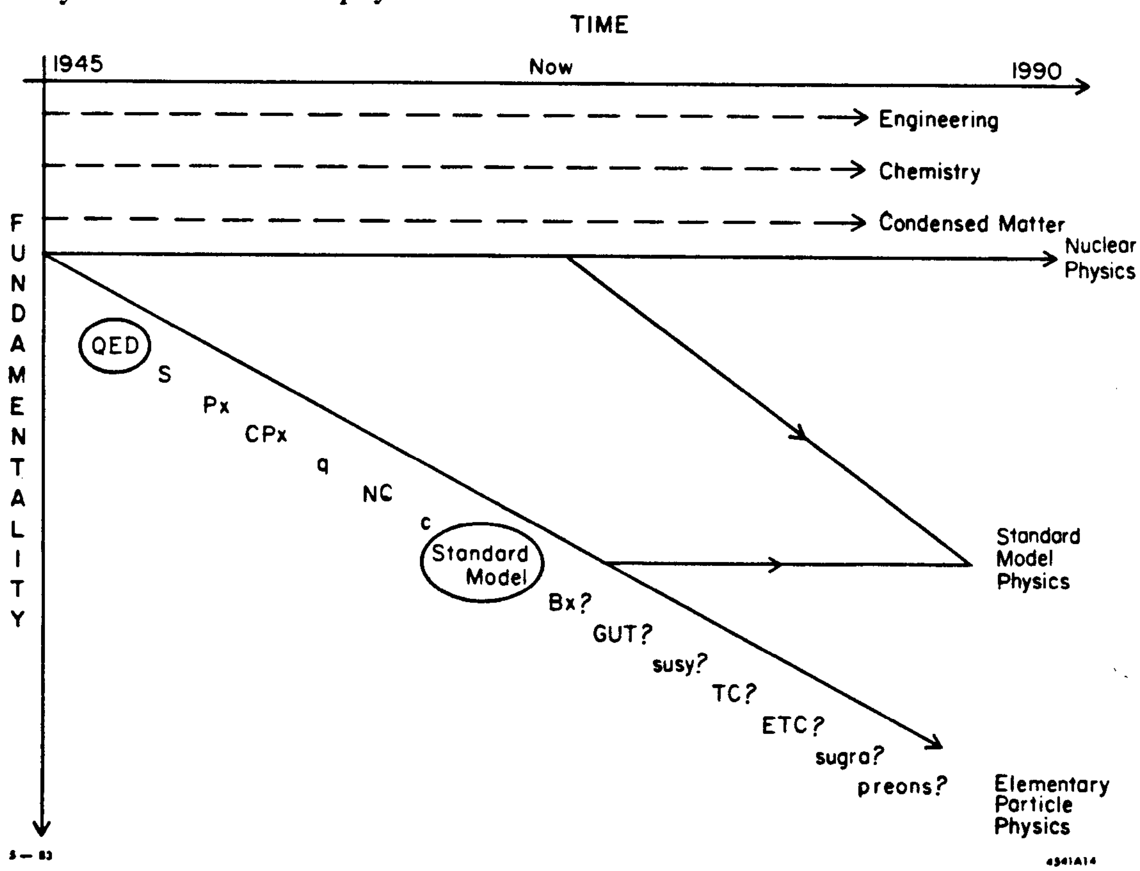


Fig. 15. An impression of the past and future history of elementary particle physics.

We have recently reached a potential parting of the ways, with the attainment of a node marking the establishment of the Standard SU(3) $\times$ SU(2) $\times$ U(1) Model. In the future one might try to develop and test the predictions of the Standard Model in greater detail, or might might look for new physical laws beyond the Standard Model. I would regard the latter as the true future province of elementary particle physics. Studying multiparticle dynamics as an end in itself would take you along the horizontal branch. This talk has been intended to indicate how multiparticle dynamics might serve as a valuable means of advancing our knowledge down the more fundamental line of elementary particle physics. Which line do you choose to follow?

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