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NEW PARTICLES AND TWO-PHOTON PHYSICS

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ABSTRACT

In a first part, I review the general theoretical arguments leading to new physics and new particles beyond the Standard Model, either in terms of supersymmetry or compositeness. Speculations about new particles expected within these schemes are then discussed in the light of recent anomalous events from the $p\bar{p}$ collider and from PETRA.
In a second part, I specifically try to evaluate the potential of $\gamma\gamma$ and $e\gamma$ collisions at PETRA/PEP and LEP energies with respect to new particle searches. Some interesting possibilities, including searches for spinless composite bosons, non-standard enhanced Higgs particles, scalar electrons ($\tilde{e}$) and $\gamma\gamma \rightarrow$'nothing' emerge.

1. New Physics beyond the Standard Model

The Standard Model of electroweak interactions [1])

$$\left[SU(2)_L \times U(1) \right]_{GSW} \tag{1}$$

is certainly in excellent health. It describes charged and neutral current reactions with ease over a wide energy range and has predicted the $W^{\pm}$ and Z^{0} vector bosons in the correct mass range. In addition to the spectacular discovery of the weak bosons [2]), evidence for the top-quark has even been announced recently [3]). Thus, only the Higgs particle is missing from the Standard Model point of view.

Let me motivate in this section why, nevertheless, a wealth of additional particles as manifestations of new physics may be expected.

The point is that besides providing a successful description of electroweak phenomena, the Standard Model also leaves open a number of important questions:
i) Which is the origin of the Fermi scale

$$\Lambda_F \equiv (\sqrt{2}\, G_F)^{-1/2} \sim 250 \text{ GeV } ? \tag{2}$$

On the one hand, Λ_F is one of the three fundamental and conspicuously different scales we know of today

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$$\Lambda_{color} \simeq 300 \text{ MeV}, \quad \Lambda_F \simeq 250 \text{ GeV} \text{ and } \Lambda_{Planck} \simeq 10^{19} \text{ GeV}. \tag{3}$$

The W, Z, Higgs, quark and lepton masses are all proportional to this single scale Λ_F of electroweak interactions. On the other hand, in the Standard Model, it arises very much ad hoc, as that value $\langle \phi \rangle$ of the scalar (Higgs) field ϕ for which the scalar potential $V(\phi)$ happens to have its absolute minimum.
ii) How can we hope to reduce the large number of free parameters (O(20)!) in the Standard Model? In a basic theory of electroweak interactions it should be possible to compute e.g. the quark and lepton masses in terms of Λ_F as well as the various KM-mixing angles.
iii) Which is the origin of the mysterious quark-lepton connection in form of families, with lepton e.m. charges being integer multiples of the quark charges?
iv) Why are there three (or more?) replicas of families, the generations?

Because of this inherent large degree of indetermination it is widely believed that the Standard Model only represents an effective theory in the wider sense; i.e. one expects that somewhere above the Fermi scale Λ_F there exists a new scale Λ_{new}, characterizing the onset of more fundamental physics. The general idea is then to relate the Fermi scale and the associated electroweak interactions to this new scale Λ_{new} and the new underlying physics.

Important clues for such a view come from the Higgs sector of the Standard Model.

Consider first the 'canonical' way of embedding electroweak interactions in schemes of grand unification (GUT), with

$$\Lambda_{new} = \Lambda_{GUT} \sim 10^{15} \text{ GeV}. \tag{4}$$

As was first realized by Wilson [4]) and subsequently emphasized by 't Hooft [5]), such a large gap

$$\Lambda_{new}/\Lambda_F \gg 1, \tag{5}$$

combined with the requirement of perturbation theory being applicable all the way up, poses the problem of an unnatural finetuning of parameters in the scalar sector. The problem originates in the quadratic divergence of radiative corrections to the Higgs mass. With $\Lambda_{new}=\Lambda_{GUT}$ acting as a cut-off the loops involving the fermions and the bosons of the theory individualy give a large contribution (Fig. 1)

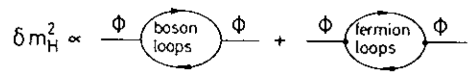

Fig. 1: Quadratically divergent radiative corrections to the Higgs mass

$$\delta m_H^2 \propto e^2 \int^{\Lambda^2_{new}} dk^2 \propto e^2 \Lambda^2_{new} \sim e^2 \Lambda^2_{GUT} \qquad (6)$$

On the other hand the validity of perturbation theory requires

$$m_H^2 + \delta m_H^2 \lesssim \mathcal{O}(\Lambda_F^2) \qquad (7)$$

Thus, a tremendous finetuning of the input parameters is required at each level of perturbation theory, unless:

A) there is a symmetry, which enforces an almost perfect cancellation of the large loop corrections

or

B) the new scale Λ_{new} is actually close to Λ_F such that possible radiative corrections of type (6) are automatically small.

These two options have found their realization in form of two popular classes of schemes of very different underlying philosophy.

(A) Supersymmetry[6] (SUSY): As is well known SUSY relates fermions and bosons with $\Delta J = 1/2$. It turns out that by making the theory supersymmetric one exactly effects the cancellation between boson and fermion loops [7], needed to avoid 'finetuning' if $\Lambda_{new}/\Lambda_F \gg 1$. The 'perturbative GUT spirit' can thus be consistently maintained. Models of supergravity [8] (SUGRA), involving SUSY in its local form, ambitiously hope to even relate the Fermi scale Λ_F directly to the 'ultimate' Planck scale

$$\Lambda_F \longleftrightarrow \Lambda_{Planck} \sim 10^{19}\ GeV. \qquad (8)$$

An important implication of these ideas is, of course, that a wealth of new particles, the SUSY partners of q, l, W, Z, γ, g, ϕ, must exist (table 1).

Table 1: Particles and their SUSY partners

J = 2	J = 3/2	J = 1	J = 1/2	J = 0
			q ⟷	$\tilde{q}$ (squark)
			l ⟷	$\tilde{l}$ (slepton)
		g ⟷	$\tilde{g}$ (gluino)	
		γ ⟷	$\tilde{\gamma}$ (photino)	
		W ⟷	$\tilde{W}$ (wino)	
		Z ⟷	$\tilde{Z}$ (zino)	
			$\tilde{\phi}_+, \tilde{\phi}_-$ (higgsino) ⟷	ϕ_+, ϕ_-
graviton ⟷	gravitino	— SUGRA ⟶	G (goldstino)	

Whereas the interactions of these new particles are essentially fixed from the requirements of gauge invariance and SUSY, their masses are quite uncertain, since SUSY cannot be an exact symmetry at low energies. However, in order not to spoil the required cancellation of divergences, [c.f. eq. (6)] the mass splitting must not be too large

$$| m^2_{particle} - m^2_{sparticle} | \lesssim \mathcal{O}(\Lambda_F^2). \qquad (9)$$

Some SUSY particles, such as the $\tilde{\gamma}$ are usually expected to be light, of the order of a few GEV, say. Intensive searches for SUSY particles are under way at PETRA/PEP and the $\bar{p}p$ collider. The familiar reaction

$$e\gamma \longrightarrow eX \qquad (10)$$

which traditionally serves to extract the photon-structure functions plays an important role in this context, as will be discussed in section 4.2.

Next, let me turn to the second possible solution of the 'finetuning' problem in the Higgs sector corresponding to $\Lambda_{new} \sim O(1 \div 5)\Lambda_F$.

(B) Compositeness:
The basic idea is that in some analogy to QCD, there exists a new gauge force ('technicolor', 'hypercolor', ...) which becomes strong in the vicinity of the Fermi scale Λ_F. Composites are formed from a set of new, strongly interacting constituents ('technifermions', 'preons', ...). Depending on how the new composite sector is thought to be interlocked with present physics, various possible scenarios emerge. Let me remind you of the essence behind the most popular schemes, proceeding in order of increasing 'radicalness'.

A crucial feature in all cases is that the Fermi scale Λ_F appears to be directly related to the confinement scale $\Lambda_{TC/HC}$ of the new underlying gauge theory.

i) Composite Higgs (Technicolor [9]): Only the Higgs scalar ϕ in the Standard Model is replaced by a boundstate

$$\phi \longrightarrow F\bar{F} \qquad (11)$$

of socalled technifermions F which, apart from the new technicolor forces, experience the same weak $[SU(2)_L \times U(1)]_{GSW}$ gauge interactions as the (elementary) quarks and leptons. The essential point is that the required spontaneous symmetry breaking in the Standard Model, is now effected dynamically through formation of technifermion condensates

$$\langle F\bar{F} \rangle \simeq \Lambda^3_{TC}, \qquad (12)$$

in analogy to the quark condensates, $\langle q\bar{q} \rangle \simeq \Lambda_C^3$, being formed in QCD during the process of confinement. Simple rescaling from QCD gives

$$\Lambda_C / f_\pi \sim \Lambda_{TC}/\Lambda_F \quad or \quad \Lambda_{TC} \sim 4\Lambda_F \sim 1\ TeV. \qquad (13)$$

The Goldstone boson analogues of the pions are subsequently 'eaten' by the W,Z bosons, whereby the latter become massive as required.

Important possible signatures of such a technicolor mechanism at much lower energies, $E \ll \Lambda_{TC}$, are extra light (pseudo) Goldstone bosons [10]. Some may even carry color (e.g. leptoquarks!).

ii) Composite quarks and leptons [11]
Here, the idea is to consider the quarks and leptons as composites of a common set of constituents ('preons'), strongly bound by the new gauge force ('hypercolor') at $E \sim \Lambda_{HC} \gtrsim \Lambda_F$. By treating quarks and leptons as boundstates one hopes to eventually compute e.g. their masses as well as to understand the origin of families and generations.

There are two drastically different ways of how the weak interactions are incorporated. The first, more conservative possibility is to view the weak interactions as fundamental gauge interactions like in the Standard Model. Spontaneous symmetry breaking then has to be effected along the lines of technicolor via a dynamical Higgs scalar being also a composite of preons. Accordingly, the hypercolor scale is as in (13)

$$\Lambda_{HC} \sim 1\ \mathrm{TeV}. \qquad (14)$$

The second possibility is much more radical [12]. The weak interactions are not associated with a fundamental gauge force and there is no Higgs mechanism at all. Instead, the weak interactions are viewed as residual hypercolor interactions among composite quarks and leptons just like the strong interactions among composite hadrons are known to be residual color interactions. The $W^\pm$ and Z^0 bosons are, correspondingly, interpreted as prominent composite vector bosons not unlike the ρ mesons in strong interactions. Since in this case

$$\Lambda_{HC} \sim \Lambda_F \simeq 250\ GeV, \qquad (15)$$

this scheme is usually termed 'nearby' compositeness. As an important experimental signature of 'nearby' compositeness one expects a possibly rich spectrum of new composite particles with masses related to the Fermi scale. There should be further fermions, both colored and uncolored (q*, l*) as well as new colored and uncolored composite bosons of various spins. When discussing 'compositeness' in relation to experiments I shall mainly refer to this scenario, simply because the new particles are supposed to be relatively 'nearby'.

Finally, let me point out that despite intensive searches no satisfactory 'Standard Composite Model' has as yet emerged. This applies both for technicolor and for composite models of quarks and leptons. Hence, 'predictions' in compositeness are mostly of qualitative nature and heavily rely on the analogy to QCD and strong interactions.

2. Anomalous Events and New Particles?

A variety of socalled anomalous events has been reported from the UA1 and UA2 detectors [13,14] at the CERN $p\bar{p}$ collider. But even at the much lower energies of PETRA (DORIS) some surprises have been claimed to exist [15,16]. In view of the meagre statistics, some of these evidences presumably are fluctuations. On the other hand, the chances are not so small, that at least some of the unusual events will be confirmed and then might provide a first glimpse on the new physics expected in the vicinity of the Fermi scale.

There has been a flood of theoretical papers trying to interpret the anomalous events as evidence for new particles from either of the two competing schemes: supersymmetry and compositeness. Let me shortly summarize the most intriguing experimental findings (as of summer 1984) along with some representative theoretical speculations.

i) Radiative Z^0 decays (UA1/UA2)
Three out of 12 Z^0 events found by UA1/UA2 involve besides a lepton pair a hard photon in addition [2]. Taken at face value this would mean a radiative rate

$$R_{UA1/UA2} \sim 20 \div 25\ \% \ , \qquad (16)$$

being in excess by about a factor 10 over the normal fraction for bremsstrahlung expected theoretically [17]. If to be interpreted in terms of new physics, these events are hard for SUSY. In contrast, there is a whole variety of suggestions from (brave) members of the compositeness camp like

(1) a new composite J = 0 partner [18] X^0 of the W,Z bosons of mass

$$m_X \simeq m(e^+e^-) \simeq 42 \div 50\ GeV\ , \qquad (17)$$

causing

$$Z^0 \to X^0\ \gamma\ . \quad X^0 \to e^+e^-, \mu^+\mu^-, \dots \qquad (18)$$

I shall come back to this suggestion in sect. 3.2, since here is a case to illustrate how crucial, complementary information can come from the two-photon channel

$$e^+e^- \to \gamma\gamma \qquad (19)$$

at much lower energies [19] (PETRA).

(2) new (composite) excited leptons [20] of mass

$$m_{\ell^*} \simeq m(\ell\gamma) \simeq 80\ GeV, \qquad (20)$$

causing

$$Z^{o} \longrightarrow \ell^{*}\,\bar{\ell} \; . \qquad \ell^{*} \hookrightarrow \ell\gamma \tag{21}$$

(3) new (composite) states X degenerate [21] with the Z° causing

$$X \longrightarrow Z^{o}\gamma \; . \qquad Z^{o} \hookrightarrow \ell\bar{\ell} \tag{22}$$

It should be noted, however, that none of these proposals seems entirely satisfactory. For instance, suggestions (1) and (2) have problems [22] to account for the 'distribution' of the $\ell^+\ell^-\gamma$ events in a Dalitz plot of $m^2(\ell^+\gamma)$, $m^2(\ell^-\gamma)$ and $m^2(\ell^+\ell^-)$.

ii) 'Monojets' and 'monophotons' (UA1)
5 + (1) 'monojets' and 1 + (1) 'monophotons' have been reported [13] by UA1:

$$\bar{p}p \longrightarrow 1 \text{ hard} \begin{Bmatrix} \text{jet} \\ '\gamma' \end{Bmatrix}_{E_T \gtrsim 35\,GeV} + (\text{jetlets}) + E_T^{miss.} \gtrsim 35\,GeV \; . \tag{23}$$

This time the SUSY advocates are much more happy. 'Monojets' can arise in SUSY schemes as signals of the production and subsequent decay of SUSY-partners such as gluinos [23] ($\tilde{g}$) or - preferably-squarks [24] ($\tilde{q}$)

$$\bar{p}p \longrightarrow \tilde{q}\,\bar{\tilde{q}} + \dots \; ; \quad \tilde{q} \longrightarrow q + \tilde{\gamma}\,(E_T^{miss}) . \tag{24}$$

The way how this process ends up fitting the experimental monojet distribution is somewhat tricky [24]. It happens through the combined effect of experimental cuts and 'jet merging' caused by the UA1 jet trigger and jet algorithms. As a result, the $\tilde{q}$ mass is claimed [24] to be in the range (Fig. 2)

$$m_{\tilde{q}} \sim 25 \div 40 \;\; GeV , \tag{25}$$

with some preference for $m_{\tilde{q}} \simeq 40$ GeV. A quite different possible SUSY-mechanism [25] for monojets and monophotons arises if SUSY happens to be spontaneously broken. In such case SUSY may be realized [26] (non-linearly) in terms of only one (massless) neutral fermion G , the goldstino being the Goldstone particle of broken SUSY *.

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* note, however, that in the popular SUGRA schemes [8] the goldstino is 'eaten' by the gravitino (Super Higgs effect) which, in turn, gets a mass $m_{3/2}$.

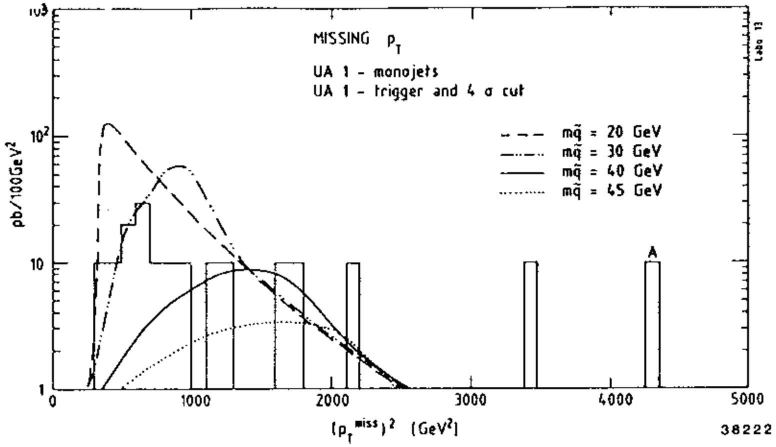

Fig. 2: Missing P_T distribution of UA1-monojets [13] compared to predictions from squark ($\tilde{q}$) production in SUSY. From J. Ellis et al. in ref. 24.

The goldstino G is pair produced [25] but not detected ($P_{\perp}^{miss}$)

$$\begin{aligned} &\bar{p}p \longrightarrow \text{jet} \; + G + G \; + \dots \qquad (\text{monojets}) \\ &\bar{p}p \longrightarrow \gamma \; + G + G \; + \dots \qquad (\text{monophotons}) . \end{aligned} \tag{26}$$

The rates and distributions can be made to roughly agree in terms of one parameter, the SUSY breaking scale. Of interest is also the resulting parameter free prediction for e^+e^- collisions [25]

$$\sigma\,(e^+e^- \longrightarrow \gamma + G + G) \simeq (\sqrt{s}/35\,GeV)^6 \cdot 10^{-2}\, pb \tag{27}$$

corresponding to $\sim$ 1 event for a typical PETRA experiment.

In sect. 3.3 I shall speculate on possibly interesting implications of such a mechanism for two-photon physics.

Nearby compositeness (technicolor) schemes also like monojets!

One proposal is to consider, in correspondence to excited leptons in (21), heavy (composite) quarks [27] ('starks')

$$gq \to q^* \to Z^0_{\hookrightarrow \nu\bar\nu}\, q \quad ; \quad m_{q^*} \sim \mathcal{O}(\Lambda_F). \qquad (28)$$

A further class of suggestions involves colored composite bosons which should naturally exist in nearby compositeness [28]. Consider first colored vector bosons (being the colored friends of W and Z) with mass

$$m_V \sim \mathcal{O}(\Lambda_F). \qquad (29)$$

(1) color octet [29] W_8, Z_8:

$$q\bar q \longrightarrow Z_8 \to Z^0_{\hookrightarrow \nu\bar\nu}\ \text{gluon}. \qquad (30)$$

They are expected in composite models of the Fritzsch-Mandelbaum type [30].

On the one hand, upon reconstruction according to (28) or (30) all monojet events indeed correspond to a mass, clustering around [31]

$$m_{Z_8(q^*)} \sim 160\ \text{GeV}. \qquad (31)$$

On the other hand, the mechanisms (28), (30) predict a specific number of corresponding events with ν's replaced by charged leptons. Apparently, they are not seen, which seems to put these proposals on somewhat shaky ground.

(2) color triplet leptoquark bosons [28,32] V_3 with Q = 2/3 decaying as

$$V_3 \Longrightarrow \begin{cases} q_{2/3} + \bar\nu \\ q_{-1/3} + \ell^+ . \end{cases} \qquad (32)$$

Monojet signals arise from the subprocess [28]

$$g\, q_{2/3} \longrightarrow V_3\ \nu_{\hookrightarrow q_{2/3}+\bar\nu,\ q_{-1/3}+\ell^+}. \qquad (33)$$

Leptoquark vector bosons of this type are expected in the class of Abbott-Farhi initiated composite models [33], characterized by the global <u>'hyper'-flavor</u> symmetry $[\alpha_c(\Lambda_F), \alpha \sim 0]$

$$SU(4)_{\text{Pati-Salam}} \supset SU(3)_{\text{Color}} \times U(1)_Y. \qquad (34)$$

It remains to be investigated in detail whether the constraints from rare decays involving $\Delta S \neq 0$ neutral currents (e.g., $K_L \to \bar\mu e$, ...) allow composite J = 1 leptoquark bosons to be sufficiently light [34,35].

As a further possible source of monojets let me finally discuss <u>spinless</u> leptoquark bosons χ of the (pseudo-) <u>Goldstone boson</u> type [36]. They arise in schemes of compositeness (technicolor) involving a spontaneous breaking of the Pati-Salam type 'hyper'-flavor symmetry (34)

$$SU(4)_{\text{Pati-Salam}} \xrightarrow{SSB} SU(3)_{\text{Color}} \times U(1)_Y. \qquad (35)$$

As their vector partners (32) they are color triplets, have charge Q = 2/3 and decay as in (32). In nearby compositeness they are expected to be light [36],

$$m_\chi \sim \sqrt{\frac{\alpha_c(\Lambda_F)}{\pi}}\ \Lambda_F \sim \mathcal{O}(40\ \text{GeV}) \ll \Lambda_F, \qquad (36)$$

(give or take a factor two).

For spinless composite leptoquarks of the Goldstone boson type there is virtually no problem with rare decays since they typically couple very weakly to quarks and leptons of one generation [36]

$$g_{\chi q \bar\ell} \sim \langle m_{q,\ell}\rangle / \Lambda_F \ll 1. \qquad (37)$$

They are, however, copiously pair-produced via color and electromagnetic gauge interactions, just like squarks ($\tilde q$)

$$\text{gluon gluon} \longrightarrow \chi\bar\chi \quad \text{in } \bar p p \text{ collisions} \qquad (38)$$

and

$$e^+e^- \longrightarrow \gamma \longrightarrow \chi\bar\chi \quad \text{in } e^+e^- \text{ collisions}. \qquad (39)$$

As to monojets, it is amusing to note, that the signals due to such leptoquarks are indistinguishable from those due to squarks [36] (eq. (24)). The production cross sections and expected mass ranges are virtually identical. Due to (37), the dominant decay of a leptoquark, associated e.g. with the second q-l generation is

$$\chi^{(2)} \longrightarrow c + \bar\nu \qquad (\gg \chi^{(2)} \longrightarrow s + \mu^+), \qquad (40)$$

which is indistinguishable from squark decay in SUSY schemes

$$\tilde q \longrightarrow q + \tilde\gamma. \qquad (41)$$

Hence, to the extent that monojets are evidence for SUSY [24] they are also evidence for leptoquarks. However, in contrast to squarks, leptoquarks also decay subdominantly into charged leptons, eq. (40), giving rise to anomalous l^+jet - l^-jet events [36,38] (Fig. 3c). This brings me to

iii) <u>Anomalous μ^+jet - μ^-jet events (CELLO/UA1)</u>

One strikingly planar μ^+jet - μ^-jet event was reported by CELLO [15] at PETRA and very recently two quite similar ones [37] by UA1 (Figs. 3a,b).

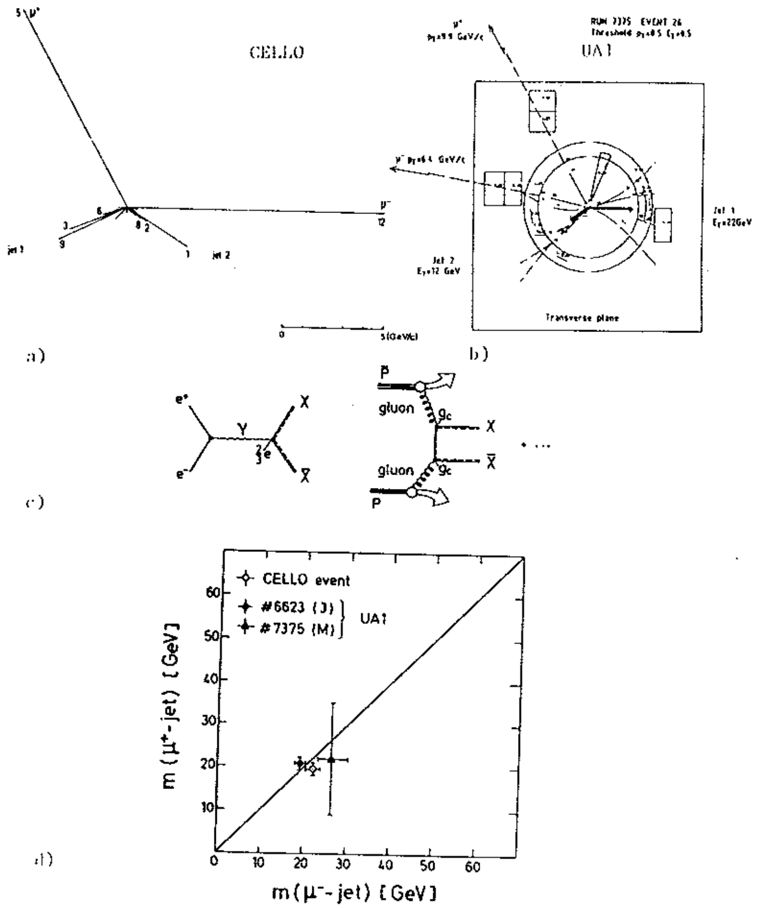

Fig. 3: Evidence for spinless leptoquark bosons χ from μ^+jet $-\mu^-$jet events
(a): CELLO event (ref.15)
(b): one of two UA1 events (ref.37)
(c): diagrams for χ pair production in $\bar{e}e$ and $\bar{p}p$ collisions
(d): consistency of $\mu^\pm$jet invariant mass with $m_\chi \simeq 20 \div 22$ GeV according to fig. 3c).

As displayed in Fig. 3d, all three events are compatible [36] with

$$m(\mu^+ jet_1) \sim m(\mu^- jet_2) \simeq 20 \div 22 \ GeV. \qquad (42)$$

Hence it is tempting to associate them with pair production [38] of the above spinless leptoquarks $\chi^{(2)}$. This means [15]

$$m_{\chi^{(2)}} \simeq 20.5 \pm 1 \ GeV. \qquad (43)$$

Such a mass is still compatible with the rough estimate (36) and also marginally within the mass range (25) needed to explain monojets along the lines of Ref. 24). The kinematics of the CELLO event ($\sqrt{s}$ = 43.45 GeV) would correspond to a $\chi\bar{\chi}$ pair being produced almost at rest. Then each χ decays into a s-μ^+ pair approximately back to back. Since two intersecting lines form a plane this mechanism naturally accounts for the almost planarity of the CELLO event [36]. Overall compatibility of the CELLO and UA1 μ^+jet - μ^-jet event rates is found for 36)

$$Br(\chi^{(2)} \to s\mu^+) \sim 0.2 \ , \qquad (44)$$

leaving

$$Br(\chi^{(2)} \to c\bar{\nu}) \sim 0.8 \qquad (45)$$

for monojet type signals and acoplanar two-jet events [in e^+e^- collisions].

iv) A narrow state $\zeta(8.3)$ in $\Upsilon(1S) \to \gamma X$ (CRYSTAL BALL)
The $\zeta(8.3)$ was announced at the Leipzig conference [16] by the CRYSTAL BALL collaboration. It had caused great excitement, e.g. as a possible candidate for a (non-standard) Higgs. Unfortunately, further runs with comparable statistics could not confirm this state [39]. I have correspondingly skipped the discussion on its implications for two-photon physics from the written version of my talk.

In conclusion, I feel it is certainly a good time to think about new particles. But, beware of inflation! Presently, more new particles have been proposed than there are 'anomalous' events. The results from the fall-run of the $\bar{p}p$ collider are certainly awaited with excitement.

3. New Particles in Photon-Photon Collisions

The problematic aspects of two-photon collisions in the context of new particle searches are quite obvious: relatively small cross sections and a strongly decreasing flux at large values of the $\gamma\gamma$ cm energy W.

As outlined in the preceeding sections, theoretical arguments lead us to expect new particles with masses somehow related to the Fermi scale Λ_F. A pessimistic guess would be

$$m_{new} > m_{W,Z} \ . \qquad (46)$$

In such a case there seems to be very little chance for two-photon physics in the foreseeable future. Even at the level of a disfavoured subprocess, e.g.

$$\text{gluon} + \text{gluon} \rightarrow \text{new particles} \tag{47}$$

the $\bar{p}p$ collider, in contrast, reaches gg energies

$$W \sim O(100\ \text{GeV}). \tag{48}$$

It may be instructive to compare the 'potential for new physics' in the gg subprocess (47) at the $\bar{p}p$ collider with that in two-photon collisions

$$\gamma + \gamma \longrightarrow \text{new particles} \tag{49}$$

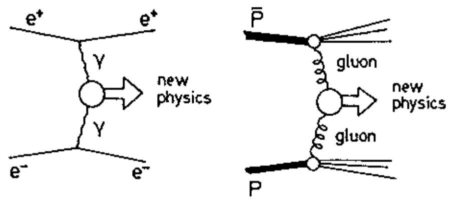

Fig. 4: $\gamma\gamma$ versus gluon gluon collisions

at PETRA/PEP and LEP in a more quantitative way (Fig. 4). For this purpose, consider the $\gamma\gamma$ and gg luminosities $d\mathcal{L}/dz$

$$\frac{d\sigma}{dz}\begin{pmatrix} e^+e^- \\ \bar{p}p \end{pmatrix} = \frac{d\mathcal{L}^{\gamma\gamma}_{gg}}{dz} \cdot \hat{\sigma}_{\substack{\gamma\gamma \\ gg}}(W), \tag{50}$$

$$\text{with} \quad z = W/\sqrt{s}, \tag{51}$$

computed via the equivalent photon approximation [40] and the gluon structure functions of Ref. 41, respectively. Since one expects

$$\hat{\sigma}_{\gamma\gamma}(W) \propto \frac{\alpha^2}{W^2} \quad \text{and} \quad \hat{\sigma}_{gg}(W) \propto \frac{\alpha_c^2}{W^2} \tag{52}$$

it seems fair to compare the quantities

$$\frac{\alpha^2}{W^2}\frac{d\mathcal{L}^{\gamma\gamma}}{dz} \quad \text{with} \quad \frac{\alpha_c^2}{W^2}\frac{d\mathcal{L}^{gg}}{dz}, \tag{53}$$

with the dimensions of a cross section as function of W. The result is

displayed in Fig. 5 and looks somewhat depressing for two-photon

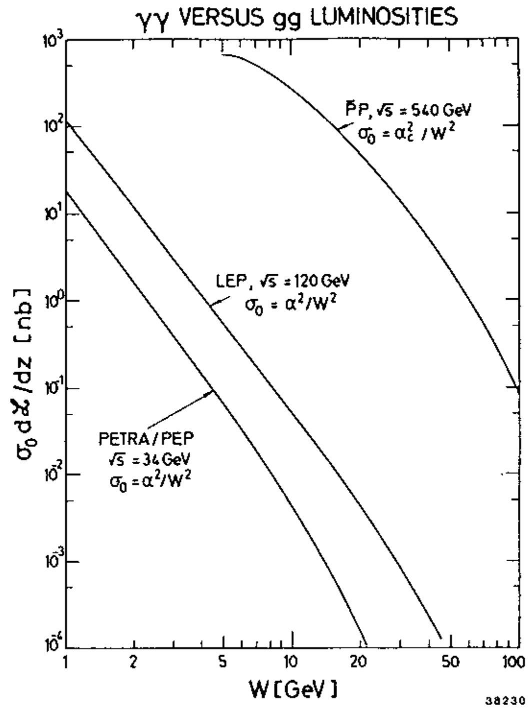

Fig. 5: Comparison of $\gamma\gamma$ and gluon gluon luminosities

physics, even at LEP.

On the other hand, both in SUSY and (nearby) compositeness it is quite possible that some of the new particles are exceptionally light

$$m \ll \Lambda_F, \tag{54}$$

and - most importantly - that their two-photon couplings are substantially enhanced. For obvious reasons (Fig. 5) I shall concentrate on such possibilities.

3.1 Enhanced Non-standard Higgs Scalars (SUSY)

First of all, it is worth recalling that the two photon (gluon) width of a relatively light Higgs scalar

$$\Gamma(\text{Higgs} \to \gamma\gamma\,[gg]), \qquad (55)$$

is an extremely interesting quantity. It involves the familiar triangle graph in Fig. 6. The point is that heavy intermediate particles excepti-

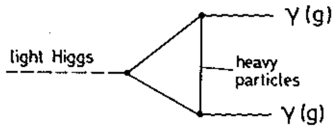

Fig. 6: Decay of a light Higgs into $\gamma\gamma$(g g)

onally do <u>not</u> decouple [42,43]. The usual loop suppression from heavies of mass m_h

$$\propto \frac{1}{m_h}, \qquad (56)$$

is compensated by the Higgs coupling

$$g_{\text{Higgs}\to\text{heavies}} = \frac{m_h}{\Lambda_F}, \qquad (57)$$

which is a consequence of spontaneous symmetry breaking. In fact [42,43]

$$\Gamma\left(\text{Higgs} \to {\gamma\gamma \atop gg}\right) \text{ 'counts' the number }^2\text{ of } \left\{ {\text{charged} \atop \text{colored}} \right\} \qquad (58)$$

heavy particles ($m_h \gg m_{\text{Higgs}}$) in the theory which is remarkable.

The problem is that for a standard GSW Higgs $\Gamma(\text{Higgs}\to\gamma\gamma)$ turns out to be hopelessly small [42,43]. The result for 3 generations is displayed in Fig. 7, to which I shall return repeatedly. Also shown are, for comparison, realistic lower limits achievable at PETRA/PEP and LEP for the quantity

$$\frac{\Gamma(R\to\gamma\gamma)\ Br(R\to X)}{M_R^3} \left[\frac{\text{keV}}{\text{GeV}^3}\right], \qquad (59)$$

as a function of the mass M_R of a given spinless particle. The region above the limits corresponds to more than 20 events at an assumed integrated luminosity

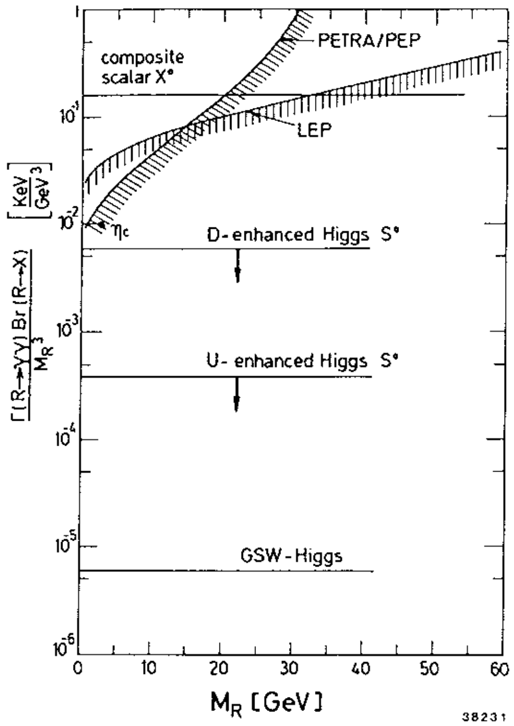

Fig. 7: Predictions for $\gamma\gamma$ widths of GSW-Higgs, non-standard U-and D-enhanced Higgs and composite scalar X^0 compared to lower limits at PETRA /PEP and LEP (see text).

$$\int \mathcal{L}\,dt = \begin{cases} 100\ \text{pb}^{-1} & \text{at PETRA/PEP;} \\ 20\ \text{pb}^{-1} & \text{at LEP; } E_b = 60\ \text{GeV} \end{cases} \qquad (60)$$

and a typical no-tag efficiency of $\varepsilon = 5\%$ for a certain final state X. Again the equivalent photon approximation was used.

The purpose of this section is to examine, how the situation may change for Higgs particles in SUSY extensions of the Standard Model.

In the minimal, non-supersymmetric case with one complex Higgs doublet ϕ, ϕ gives mass to up quarks, whereas ϕ^* gives mass to down quarks and charged leptons via spontaneous symmetry breaking $\langle\phi\rangle \neq 0$. In SUSY, however, ϕ and ϕ^* cannot be both members of SUSY multiplets and hence at least <u>two</u> (complex) scalar doublets ϕ_+ and ϕ_- must exist [6,44]

$$\phi \to \phi_+ \quad ; \quad \phi^* \to \phi_- \qquad (61)$$

Of these 2 x 2 x 2 = 8 real scalar degrees of freedom, three again serve to give mass to $W^{\pm}, Z^{o}$. Therefore, in SUSY extensions instead of one neutral, physical Higgs scalar H^{o}, there are now (at least) five physical Higgses, including two charged ones!

$$H^{o}, \quad S^{o}, P^{o}, \quad S^{\pm}. \tag{62}$$

The Yukawa couplings of these Higgses to fermions now involve two additional parameters, a H^{o}-S^{o} mixing angle Θ as well as the ratio of the two vacuum expectation values

$$0 < \frac{\langle\phi_+\rangle}{\langle\phi_-\rangle} \equiv tg\alpha < \infty . \tag{63}$$

Since $tg\alpha$, a priori may have any value, the Higgs phenomenology is, correspondingly, much richer. In particular, dramatic enhancements of certain Yukawa couplings may arise if

$$tg\alpha \ll 1 \qquad \text{or} \qquad tg\alpha \gg 1 . \tag{64}$$

Specifically, one finds for the three neutral Higgses, H^{o}, S^{o} and P^{o}, contributing via the triangle (Fig. 6) to $\gamma\gamma$, the following general couplings [44] to up quarks U = (u,c,t), down quarks D = (d,s,b) and charged leptons $L^{-} = (e,\mu,\tau)$

$$\mathcal{L}_{Yuk.} = \sum_{f=U,D,L^-} \frac{m_f}{\Lambda_F} \cdot X_f \left\{ f_L^{+} f_R \cdot Higgs + h.c. \right\} . \tag{65}$$

Of central interest in (65) are, of course, the 'enhancement' factors X_{U,D,L^-} which, as functions of the parameters α and Θ are displayed in table 2. Note, in particular, that $X_{L^-} = X_D \gg 1$ implies $X_U < 1$ and vice versa.

Table 2 Enhancement factors for Yukawa couplings of neutral Higgses in minimal SUSY extensions, compared to the single Higgs case (H^{o}) in the Standard Model (GSW)

SUSY-GSW	$X_{L^-} = X_D$	X_U
H^{o}	$\cos\Theta / \cos\alpha$	$\sin\Theta / \sin\alpha$
S^{o}	$\sin\Theta / \cos\alpha$	$-\cos\Theta / \sin\alpha$
P^{o}	$-i \sin\alpha / \cos\alpha$	$-i \cos\alpha / \sin\alpha$
GSW		
H^{o}	1	1

In addition, the masses of the Higgs particles also depend on $tg\alpha = \langle\phi_+\rangle / \langle\phi_-\rangle$ and they are related among each other [44]. For large classes of (broken) SUSY schemes one finds in case of enhancement of either the U-sector ($X_U \gg 1$) or the D-sector ($X_D = X_{L^-} \gg 1$)

$$m_{S^o} \sim m_{P^o} \text{ arbitrary, but probably light} \tag{66}$$

and

$$m_{H^o} > m_{Z^o}, \quad m_{S^{\pm}} > m_W . \tag{67}$$

Note that the mass of (S^{o},P^{o}) is not even bounded from below by the Weinberg-Linde bound [45]

$$m_{Higgs} \gtrsim 10 \text{ GeV} , \tag{68}$$

since it only applies to the single Higgs case!

Hence, it seems, there is a characteristic signature of SUSY extensions of the Standard Model, worth being looked for in $\gamma\gamma$ collisions. The (S^{o},P^{o}) Higgses could both be possibly light and strongly enhanced in $\gamma\gamma$ via the fermions in the triangle (Fig. 6). For instance

$$\frac{\Gamma(S^{o}\to\gamma\gamma)}{m_{S^o}^{3}} \approx (GSW)_{\substack{fermionic\\ contribution}} \cdot X_f^{2} . \tag{69}$$

Furthermore, in SUSY, all the charged (heavy) SUSY partners ($\tilde{q}$, $\tilde{l}$, etc) contribute via the triangle. According to (58), a measurement of $\Gamma(S^{o},P^{o}\to\gamma\gamma)$ could thus in addition provide information on the existence of SUSY degrees of freedom, too heavy to be accessible directly!

An enhancement of the type

$$X_{L^-} = X_D \gg 1 \tag{70}$$

would be most favorable for two-photon collisions since it implies a large branching ratio of (S^{o},P^{o}) into charged leptons and thus a good detection efficiently, $\varepsilon \gg 5\%$: On the other hand, (70) requires

$$m_{S^o,P^o} > m_{\Upsilon} \simeq 9.46 \text{ GeV}, \tag{71}$$

else (S^{o},P^{o}) would probably have been seen in

$$\Upsilon \longrightarrow \gamma X \tag{72}$$

(see Fig. 8). The $\zeta(8.3)$, which was unfortunately not confirmed [16,39] (see sect. 2), could have been, of course, an exciting application for the type of exercise presented in this section.

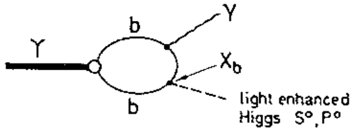

Fig. 8: A light D-enhanced Higgs in Υ decay

Finally, as a warning against too much optimism, let me point out that the crucial enhancement factors X_f cannot be too large if one insists on perturbation theory to remain valid. Then, in analogy to the perturbative unitarity bound [46] for the GSW Higgs mass

$$m_H \lesssim \sqrt{8\pi}\cdot\Lambda_F \simeq 1.2\ \mathrm{TeV} \tag{73}$$

from

$$W_L \bar{W}_L \rightarrow (\mathrm{Higgs}) \rightarrow W_L \bar{W}_L\ , \tag{74}$$

there is a corresponding limit [47]

$$m_f \cdot X_f \lesssim \sqrt{2\pi}\cdot\Lambda_F \simeq 640\ \mathrm{GeV} \tag{75}$$

from

$$f\bar{f} \longrightarrow (\mathrm{Higgs}) \rightarrow f\bar{f},\ W\bar{W}. \tag{76}$$

Using $m_b \simeq 5$ GeV and $m_t \simeq 42$ GeV, one finds

$$X_{L^-} = X_D \lesssim 130 \quad \mathrm{and} \quad X_U \lesssim 15\ . \tag{77}$$

The corresponding upper limits for $\frac{\Gamma_{S^o\to\gamma\gamma}}{m_S}$ from (77) (ignoring contributions from SUSY partners) are displayed again in Fig. 7. Unfortunately they are still small. It should be emphasized, however, that the perturbative bounds (73), (75) are not sacred. In fact, related to the anomalous pp collider events the possibility of a strongly interacting Higgs sector (violating (73) and/or (75)) has received much attention recently [48].

3.2 Enhanced Scalars X° in Nearby Compositeness

In this section, I want to consider 'nearby' compositeness, i.e. the weak W and Z° bosons are viewed as prominent composite spin 1 bosons along with composite quarks and leptons [12]. The issue is then to look for possible composite J = 0 partners X° of the W,Z vector bosons. By playing the role of J = 0 'ground states', they might well be lighter than their J = 1 counter parts W and Z°. As a basis for the discussion, let us consider a tentative mass range

$$30 \div 50\ \mathrm{GeV} \lesssim m_{X^o} < m_W \tag{78}$$

If such X° bosons were to exist in form of (pseudo) Goldstone bosons (in analogy to π^o) they could of course be even lighter.

The essential point I want to make is that the $\gamma\gamma$ and $Z^o\gamma$ channels offer some unique possibilities to detect such bosons - if they exist - even though they are a priori not expected to be very light.

(1) If X° is a member of a weak isospin triplet like $W^\pm, Z^o$, it does not couple to two (or more) gluons since those are iso-singlets.

(2) If X° is an isospin singlet its coupling to two gluons is still very weak if its constituents carry no color, like it is the case for $W^\pm, Z^o$ in a variety of popular models.

(3) The coupling of X° to light quarks and leptons is probably strongly suppressed

$$g_{X^o f\bar{f}} \propto \frac{m_f}{\Lambda_F} \ll 1\ , \tag{79}$$

for reasons of chiral symmetry [49,50].

The two photon channel then remains as a dominant [50] decay mode of X°!

Next I want to argue that the $X^o\gamma\gamma$ and $Z^oX^o\gamma$ couplings are strongly enhanced for a composite X° particle as compared to an elementary scalar [51]. The crucial difference becomes qualitatively obvious from Fig. 9. Whereas an elementary Higgs-type scalar couples to $\gamma\gamma$ and $Z^o\gamma$

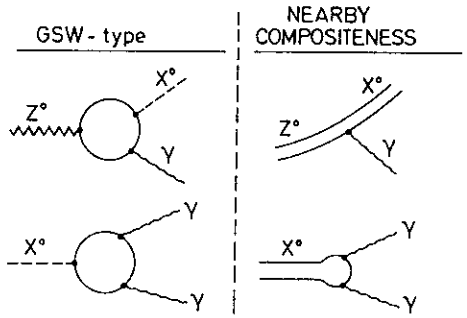

Fig. 9: $X^o\gamma\gamma$ and $Z^oX^o\gamma$ couplings for a composite scalar X^o and an elementary, Higgs-type scalar X^o, respectively

only via strongly suppressed loops (ignoring possible enhancements à la sect. 3.1), there is a direct coupling of photons to the common charged constituents of X° and Z° in case of compositeness! In order to obtain estimates of these couplings one may employ either methods familiar from onium-bound states [51] or the concept of W-dominance [52] (Fig. 10) which I prefer. The results happen to be quite similar, though.

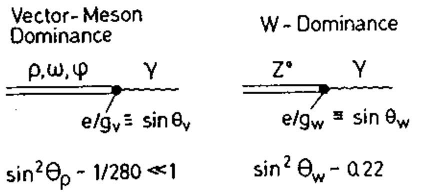

Fig. 10: Analogy of W- dominance to vector - meson dominance

W-dominance is pictured in direct analogy to vector-meson dominance [53] of the electromagnetic current in strong interactions (Fig. 10). The corresponding direct coupling of the photon to the composite Z^o is

$$e/g_W = \sin\Theta_W \simeq \sqrt{0.22} \qquad (80)$$

It is comforting that the requirement of W-dominance in its stronger, operator form [28] ('current-field identity')

$$J_\mu^{e.m} = \frac{m_W^2}{g_W} \cdot W_\mu^3 + J_\mu^\gamma \qquad (81)$$

gives very sensible results as to the effective weak interactions of composite $\vec{W}$'s, quarks and leptons. Following and generalizing the classical, analogous construction by Lee and Zumino [54] for composite ρ's and nucleons one finds a result [28] closely mimicking the standard GSW model for $E \lesssim m_W$!

$$\mathcal{L}_{eff}^{weak}(\vec{W}, q, \ell, \gamma) = \mathcal{L}_{GSW}^{\text{unitary gauge}} \ \text{without physical Higgs} + \dots \qquad (82)$$

Now, let us estimate $\Gamma(X^o \to \gamma\gamma)$. As depicted in Fig. 11, W-dominance

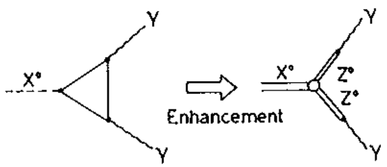

Fig. 11: Enhancement of the $X^o\gamma\gamma$ coupling in 'nearby' compositeness

relates $g_{X^o\gamma\gamma}$ to an effective X^oZZ coupling among composites for which one may naturally assume

$$g_{X^oZZ} \sim g_{Zf\bar{f}} = g_{ZWW} = g_W \simeq 0.64. \qquad (83)$$

One obtains a very large width

$$\frac{\Gamma(X^o \to \gamma\gamma)}{m_{X^o}^3} \simeq \frac{4}{3}\,\frac{\Gamma(Z \to e^+e^-)}{m_Z^3} \simeq 0.16\ \frac{keV}{GeV^3}. \qquad (84)$$

As becomes apparent from Fig. 7, there seems indeed to be a good chance for two-photon physics, in particular at LEP, up to masses

$$m_{X^o} \sim 50 \text{ GeV}. \qquad (85)$$

Even, if nothing is found, the resulting limits could be of great importance. Here is an instructive illustration, involving experiments at PETRA in particular also the two photon channel

$$e^+e^- \longrightarrow \gamma\gamma. \qquad (86)$$

We return to the proposal mentioned in sect. 2, that the anomalous radiative Z^o decays

$$Z^o \longrightarrow e^+e^- \gamma_{hard}, \qquad (87)$$

seen by UA1/UA2 are mediated by a spinless composite boson [18] X^o (Fig. 12) of mass $m_X \simeq m_{e^+e^-} \simeq 45\text{-}50$ GeV.

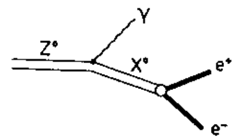

Fig. 12: A composite scalar X^o mediating radiative Z^o decays

The UA1/UA2 rate (16) then implies

$$\Gamma(Z^o \to e^+e^-\gamma) = \varepsilon\,\frac{\Gamma(Z^o \to X^o\gamma)\,\Gamma(X^o \to e^+e^-)}{\Gamma_{X^o}} \sim \mathcal{O}(15\,MeV) \qquad (88)$$

In eq. (88) $\varepsilon = 1$ applies for a single spinless boson and $\varepsilon = 2$ for a (degenerate) parity doublet [19].
W-dominance gives [19] (Fig. 13)

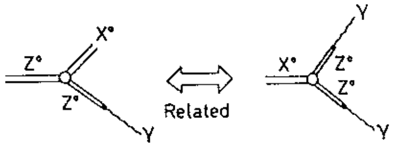

Fig. 13: Relation of $Z^oX^o\gamma$ and $X^o\gamma\gamma$ couplings from W- dominance

$$\frac{\Gamma(Z^o\to X^o\gamma)}{\Gamma(X^o\to\gamma\gamma)} = \frac{2}{3}\left(\frac{m_Z^2 - m_X^2}{m_Z m_X}\right)^3 \frac{1}{\sin^2\Theta_W} \equiv r \sim \mathcal{O}(10). \qquad (89)$$

Thus, the radiative Z^o decay rate directly constrains searches for the X^o boson in the $e^+e^-\to\gamma\gamma$ channel via the relation following from eqs. (88), (89)

$$\frac{\varepsilon\,\Gamma(X^o\to\gamma\gamma)\,\Gamma(X^o\to e^+e^-)}{\Gamma_{X^o}} \sim \frac{15\ \text{MeV}}{r} \simeq 1.5\ \text{MeV}. \qquad (90)$$

A search was performed by all PETRA groups[55-58], but no comparable signal was seen (Fig. 14a), implying

$$m_{X^o} > 46.8\ \text{GeV}. \qquad (91)$$

Since $\Gamma(X^o\to\gamma\gamma)/\Gamma_{X^o} \leq 1$, eq. (90), moreover provides the lower bound[19]

$$\varepsilon\ \Gamma(X^o\to e^+e^-) \gtrsim 1.5\ \text{MeV}, \qquad (92)$$

involving only W-dominance and the UA1/UA2 results. It turned out to be in strong contradiction to the upper bounds from all PETRA groups[55-58] for $M_{X^o} \lesssim 46.8$ GeV (Fig. 14b). In view of the bound (92), Bhabha scattering,

$$e^+e^- \longrightarrow e^+e^- \qquad (93)$$

is also very restrictive (Fig. 14c).

a)

b)

c)

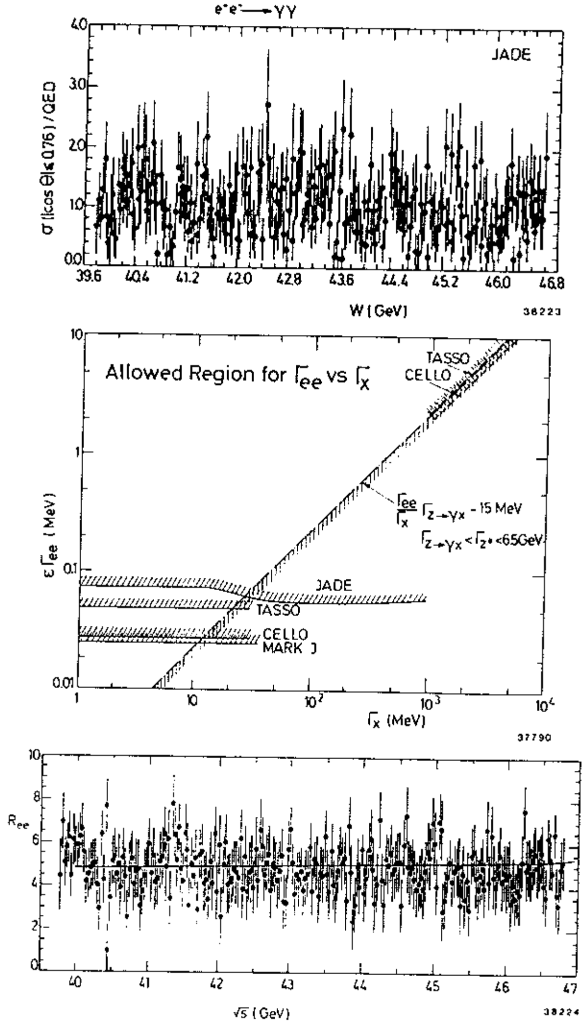

Fig. 14: Typical results of PETRA searches for a composite scalar X^o.
(a): $e^+e^-\to\gamma\gamma$ data from ref.55,
(b): summary of PETRA bounds for $\Gamma(X^o\to e^+e^-)$ vs. Γ_{X^o} (ref.55),
(c): $e^+e^-\to e^+e^-$ data from ref.56.

For X^o boson masses above the PETRA energy range a sizeable contribution is still expected in reaction (93) due to γ-X^o interference[19]. The solid line in Fig. 14c displays this contribution for

$$m_{X^o} = 49\ \text{GeV},\quad \varepsilon\,\Gamma(X^o \to e^+e^-) = 4\ \text{MeV},\quad \Gamma_{X^o} \lesssim 500\ \text{MeV}. \qquad (94)$$

Such an X^o boson is ruled out at the 95% confidence level.

Alltogether, the experiments performed at PETRA have provided strong evidence against an interpretation of the observed radiative Z^o decays in terms of a composite X^o boson. First of all these results, certainly, represent an encouraging illustration that crucial information on new physics can also come from a two-photon channel at comparatively low energies. Furthermore, it is important to realize that these results do not at all rule out the existence of a spinless composite X^o boson with

$$m_{X^o} \lesssim 50\ \text{GeV}. \qquad (95)$$

On the contrary, they merely tie in with the mentioned theoretical arguments[49,50] which suggest an extremely weak coupling of X^o to light fermions as a consequence of chiral symmetry (c.f. eqs. (73) and (92)). Thus, in view of these results and the expected large coupling of X^o to two photons, eq. (84), it appears worthwhile to look for it in $\gamma\gamma$ collisions (LEP).

3.3 $\gamma\gamma \to$ 'nothing'?

The content of this short section is admittedly very speculative, but nevertheless potentially interesting. I just would like to stimulate some thinking about the 'impossible' reaction

$$\gamma\gamma \longrightarrow \text{'nothing'}. \qquad (96)$$

Here are some reflections on two obvious questions

i) Why could reaction (96) be interesting?
ii) Is there a chance to isolate it experimentally?

(i): The only conventional source, $\gamma\gamma \to \nu\bar{\nu}$, is extremely small, since ν's carry no electric charge. What about possible unconventional sources? Here, SUSY comes to mind, since a priori there are various possibilities to turn two photons into a pair of ($\sim$ massless) invisible SUSY particles:

$$\gamma\gamma \to GG \quad \text{(goldstinos [gravitinos])} \qquad (97)$$
$$\gamma\gamma \to \tilde{\gamma}\tilde{\gamma} \quad \text{(photinos)}. \qquad (98)$$

As discussed in sect. 2, goldstino pair production within a nonlinear realization[26] of (spontaneously broken) SUSY has been suggested in Ref. 25 as a possible explanation of the monojets and monophotons observed by UA1 at the $\bar{p}p$ collider (c.f. eqs. (26)). In this scheme one may even compute the rate (97) without free parameters, which seems worthwhile. On general grounds (low-energy theorem) one expects a very strong dependence of $\sigma(\gamma\gamma \to GG)$ on the $\gamma\gamma$ energy W, analogously to eq. (27).

(ii): If possible at all, an experimental isolation of reaction (96), is certainly going to be difficult. In e^+e^- collisions, it corresponds to events of the type

$$e^+e^- \longrightarrow e^+e^- + \text{'nothing'}. \qquad (99)$$

Double tagging at large angles should at least strongly reduce the probability that photons radiated from the scattered electrons are lost in the beam pipe and fake reaction (99). However, selectron ($\tilde{e}$) pair production with subsequent decay

$$\tilde{e} \longrightarrow e + \tilde{\gamma}, \qquad (100)$$

represents another possible SUSY 'background' to relation (96) with a similar signature (99). More promising seems to be an attempt to extract reaction (96) from direct $\gamma e^\pm$ collisions involving real photons. As has been discussed in the literature[59], $\gamma_{real}e^\pm$ collisions might be realized at the Stanford Linear Collider (SLC) with γe energies $\lesssim$ 90 GeV and good luminosity by means of a laser beam.

4 New Particles in $e^\pm\gamma$ Collisions

4.1 Photon Structure Functions

Deep inelastic scattering of electrons on a (quasi-real) photon 'target' represents the standard process from which the photon structure functions $F_2^\gamma(x,Q^2)$ and $F_L^\gamma(x,Q^2)$ are extracted.

Modifications of the familiar QCD predictions for F_2^γ and F_L^γ will arise, if new particles exist.

For the case of SUSY, the effects due to squark ($\tilde{q}$) and gluino ($\tilde{g}$) production have been calculated explicitely for the structure functions of the photon[60-62] (LEP), [as well as of the nucleon[63] (HERA)] . In general, of course, the results strongly depend on the masses of the new particles under consideration.

Let me remind you here of a characteristic effect related to their spins[61,62]. The point is that the longitudinal structure function of the photon $F_L^\gamma(x,Q^2)$ turns out to be quite sensitive to the presence of spinless particles such as

	squarks $\tilde{q}$	in SUSY (Fig. 15)	
or			(101)
	leptoquarks χ X^o bosons ⋮	in 'nearby' compositeness . (c.f. sects. 2 and 3.2)	

It is well known that in QCD, with only J = 1/2 quarks and J = 1 gluons, F_L^γ exhibits no scaling violations both at the parton and $O(\alpha_S)$ levels for large $Q^2 \gg m_q^2$

$$F_L^\gamma(x, Q^2) = \frac{\alpha}{2\pi} f(x) . \qquad (102)$$

Transient Q^2 dependences will, of course, arise if new, heavy quark thresholds are crossed, e.g. around

$$Q^2 \approx 4\, m_{top}^2 . \qquad (103)$$

In contrast, new J = 0 particles give rise to an asymptotically surviving, qualitatively different behaviour[61,62)]

$$F_L^\gamma /_{J=0} \propto \log Q^2 , \qquad (104)$$

as a consequence of simple helicity conservation arguments (Fig. 15).

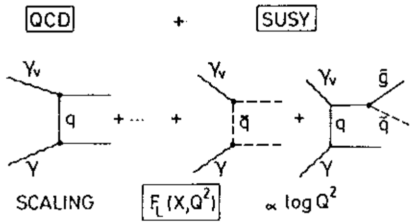

Fig. 15: Spinless squarks ($\tilde{q}$) causing a log Q^2 scaling violation in the longitudinal photon structure function $F_L(x,Q^2)$

However, in order to be able to explore this discriminating effect much effort will have to go into an experimental isolation of F_L^γ at LEP energies.

4.2 Hunting for Selectrons ($\tilde{e}$)

A detailed computation of cross sections with general couplings [and polarization states] for exclusive $e^\pm\gamma$ collisions

$$e^\pm\gamma \rightarrow \text{boson} + \text{fermion} \qquad (105)$$

was performed in ref. 64.

An exclusive SUSY process of special interest is (Fig. 16)

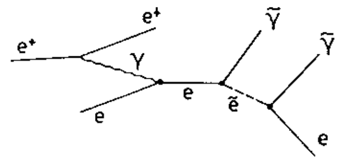

Fig. 16: Single selectron ($\tilde{e}$) production in $e\gamma$ scattering

$$e\gamma \longrightarrow \tilde{e}\,\tilde{\gamma} \quad \hookrightarrow e\tilde{\gamma} \qquad (106)$$

In contrast to pair production of selectrons ($\tilde{e}$) in e^+e^- collisions (Fig. 17)

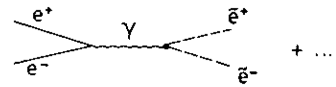

Fig. 17: Selectron ($\tilde{e}$) pair production

$$e^+e^- \longrightarrow \tilde{e}^+\tilde{e}^- \quad \hookrightarrow e^+e^- + \text{'nothing'} , \qquad (107)$$

it offers an opportunity to search for selectrons with mass higher than the e^+e^- beam energy. The $e^\pm$ which radiates the quasi-real photon in (106) goes mostly along the beam direction, and is not observed. Thus, the final state to be observed in (106) consists of a single electron coming from the $\tilde{e} \rightarrow e + \tilde{\gamma}$ decay , with no other visible particles. The cross section for this process (106) was first calculated in ref. 65 for a massless photino. A generalization to massive photinos may be found in ref. 66. Fig. 18, taken from JADE[67)] illustrates the type of bounds on selectron and photino masses one may achieve at PETRA. Note, the selectron mass bound

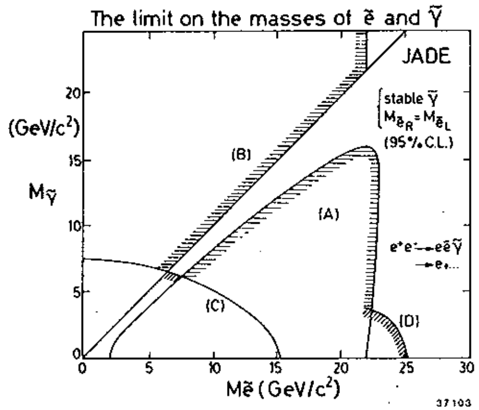

Fig. 18: Lower limits on selectron ($\tilde{e}$) and photino ($\tilde{\gamma}$) masses (ref.67)
(A): from $e^+e^- \to \tilde{e}^+\tilde{e}^-$ with $m(\tilde{e}) > m(\tilde{\gamma})$,
(B): from $e^+e^- \to \tilde{e}^+\tilde{e}^-$ with stable $\tilde{e}$, $m(\tilde{e}) < m(\tilde{\gamma})$,
(C): from $e^+e^- \to \gamma\tilde{\gamma}\tilde{\gamma}$
(D): from $e\gamma \to \tilde{e}\tilde{\gamma}$, $\tilde{e} \to e + \tilde{\gamma}$

$$m_{\tilde{e}} \gtrsim 25.2 \text{ GeV} \quad \text{for } m_{\tilde{\gamma}} = 0, \tag{108}$$

essentially coming from the $e\gamma$ process (106). For comparison, the MARK-II and MAC collaborations obtained lower $\tilde{e}$ mass limits of 22.2 and 22.4 GeV, respectively[68].

Acknowledgements

I wish to thank Hermann Kolanoski for helpful discussions on experimental aspects of two-photon physics, and Richard Lander and colleagues for organizing an enjoyable and fruitful meeting.

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