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ABSTRACT

We survey the possibilities which high energy electron-proton collisions in the HERA energy range provide to search for new particles and to test for new interactions. Particular emphasis is placed on the production of leptoquarks and supersymmetric particles. Helicity changing contact interactions are briefly discussed.

*) Lectures given at the XV International Winter Meeting on Fundamental Physics, Sevilla, Febrero 1987

1) Introduction

At the storage ring HERA currently under construction electron-proton collisions will be observed with center-of-mass energies up to $\sqrt{s}$=314 GeV and momentum transfers up to $Q^2 \approx 10^5$ GeV^2. Hence it will be possible to produce new elementary particles with masses up to the maximal center-of-mass energy and to test for new interactions down to distances $r_{min} \approx (200\ \text{GeV})^{-1} \approx 10^{-16}$ cm.

Since the proton is composed out of quarks, antiquarks and gluons with average momenta given by the inverse size of the proton, i.e., $\langle k_p \rangle \approx \frac{1}{\langle r_p \rangle} \approx 300$ MeV, high energy electron-proton collisions are scattering processes between the following quasi-free partons:

electron beam: e, γ

proton beam: $u, d, g, \gamma, \bar{u}, \bar{d}, s, \bar{s}, \ldots,$

where we have included the photons from electron and quark Bremsstrahlung.

The highest masses of new particles are accessible in single particle resonance production. In Sec. 2 we give a list of possible processes and discuss in detail the two most important cases: excited quarks and leptons, and leptoquarks. If the new particles carry a new conserved quantum number they can only be produced in pairs. This is the case in most supersymmetric models where a discrete symmetry, R-parity, distinguishes between ordinary quarks, leptons and gauge bosons and their predicted superpartners. Pair production of supersymmetric particles is discussed in Sec. 3. In addition to the most important process, the pair production of scalar leptons and quarks, we consider the production of goldstino pairs in models with nonlinearly realized supersymmetry. At low energies virtual effects of new particles can be discribed by local, nonrenormalizable contact interactions. In Sec. 4 we discuss chirally invariant, helicity changing four-fermion interactions, in particular their effect on the Callan-Gross relation.

We conclude with a list of particle masses, and interaction scales which will be accessible at HERA. Much of the material discussed in these lectures is based on previous review articles (cf. [1]-[4]) where further references can be found.

2) Resonance production of new particles

The various new particles, which can be produced in two-parton processes in electron-proton collisions, are shown in Figs. (1a)-(1g). The corresponding interaction lagrangians read

$$L_1 = \lambda_1 \bar{u}^c e\, S + \text{h.c.}, \qquad (1a)$$

$$L_2 = \lambda_2 \bar{u} e\, R + \text{h.c.}, \qquad (1b)$$

$$L_3 = \frac{\lambda_3}{\Lambda} \bar{u}^* \sigma^{\mu\nu} u\, F_{\mu\nu} + \text{h.c.}, \qquad (1c)$$

$$L_4 = \frac{\lambda_4}{\Lambda} \bar{e}^* \sigma^{\mu\nu} e\, F_{\mu\nu} + \text{h.c.}, \qquad (1d)$$

$$L_5 = \frac{\lambda_5}{\Lambda} \bar{e}_8^A \sigma^{\mu\nu} e\, G^A_{\mu\nu} + \text{h.c.}, \qquad (1e)$$

$$L_6 = \frac{\lambda_6}{\Lambda} G^A_{\mu\nu} F^{\mu\nu}\, \chi_8^A + \text{h.c.}, \qquad (1f)$$

$$L_7 = \frac{\lambda_7}{\Lambda} F_{\mu\nu} F^{\mu\nu}\, \chi + \text{h.c.}\ . \qquad (1g)$$

Here S and R are scalar leptoquarks, u^* (e^*) denotes an "excited", spin-$\frac{1}{2}$ u-quark (electron), $F_{\mu\nu}$ and $G^A_{\mu\nu}$ are electric and colour field strengths, e_8 is a colour-oktet electron ("leptogluon"), and χ and χ_8 are colour-singlet and colour-oktet scalars. Λ is a mass scale which has to be O(1 TeV) if the listed interactions are to be of interest for HERA. One can easily write down similar interaction terms involving other quark species, dual field strength tensors,

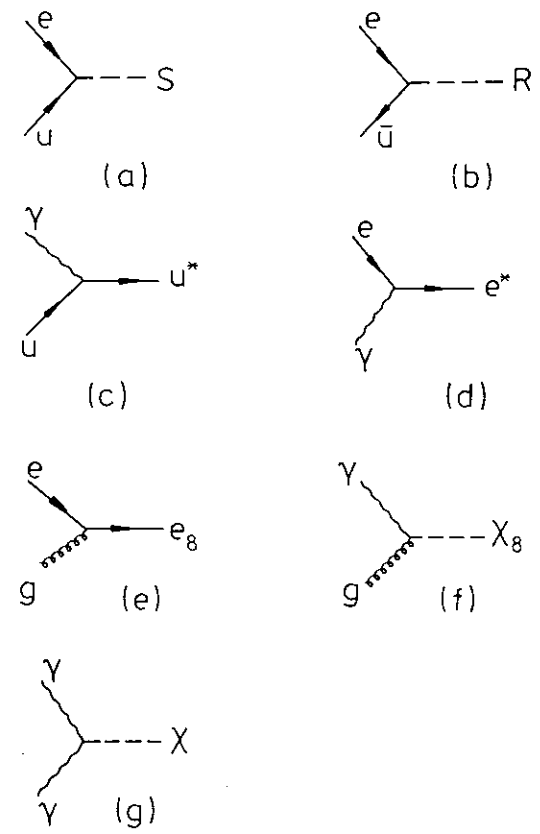

Fig. 1 Production of new particles in two-parton processes in ep-collisions: leptoquarks (a,b), excited quarks and leptons (c,d), leptogluon (e), colour-octet and colour-singlet scalars (f,g)

$W^\pm$- and Z-bosons, and further new particles with higher spins, whose couplings are generally proportional to higher powers of $\frac{1}{\Lambda}$. The off-diagonal couplings of photon and gluon have to be of magnetic moment type, since the electromagnetic and colour currents are conserved. An important aspect of eqs. (1) is that the only new particles with renormalizable couplings to the available partons are leptoquarks, whose existence is predicted in almost all extensions of the standard electroweak theory. Nonrenormalizable interactions proportional to $\frac{1}{\Lambda}$ occur in theories of quark-lepton substructure as well as through radiative corrections in the standard model. However, in the latter case the coefficients $\frac{\lambda_i}{\Lambda}$ are much smaller than the electroweak scale $\sqrt{G_F} \approx (300\ \mathrm{GeV})^{-1}$. For instance, the two-photon coupling of the standard model Higgs boson H^0 is given by eq. (1g) $(x=H^0)$ with [5]

$$\frac{\lambda_7}{\Lambda} = \frac{\alpha}{8\pi}\,(\sqrt{2}G_F)^{\frac{1}{2}}\,\left(-7 + \frac{4}{3}\sum_f Q_f^2\right), \tag{2}$$

where Q_f is the charge of a "heavy" fermion $(m_f/m_{H^0} > 0.2)$. Hence, if quarks and leptons have no substructure at the Fermi scale, the only resonances which one may hope to find in electron-proton collisions below 1 TeV are leptoquarks.

The production cross sections for resonances in the various two-parton channels are easily estimated. For a narrow resonance $(\Gamma_R \ll M_R)$ one has

$$\sigma_R \propto \delta(\hat{s}-M_R^2)$$

$$= \frac{1}{s}\,\delta(x-x_R), \tag{3a}$$

$$\hat{s} = xs,\ 0 \leq x \leq 1,\ x_R = \frac{M_R^2}{s}\ , \tag{3b}$$

where $\sqrt{\hat{s}}$ is the parton center-of-mass energy (cf. Fig. 2).

The proportionality constant is easily obtained from the Breit-Wigner resonance cross section

$$\sigma_R(s) = \frac{(2J_R+1)}{h_1h_2}\,\frac{4\pi}{\hat{s}}\qquad \frac{\Gamma_{in}\Gamma_R}{(\sqrt{\hat{s}}-M_R)^2 + \frac{\Gamma_R^2}{4}}\ , \tag{4}$$

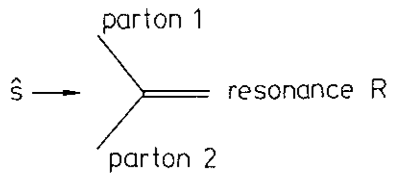

Fig. 2 Resonance production in a two-parton subprocess

which becomes in the narrow-width approximation

$$\sigma_R(s) = \frac{(2J_R+1)}{h_1h_2}\,16\pi^2\,\frac{\Gamma_{in}}{M_R}\,\frac{1}{s}\,\delta(x-x_R). \tag{5}$$

Here J_R and M_R are spin and mass of the resonance $(x_R=\frac{M_R^2}{s})$, Γ_{in} is the partial decay width into the initial state, and h_1, h_2 are the helicity multiplicities of the two partons (e.g., $h(e_L) = 1$, $h(\gamma) = 2$, etc.). With

$$\frac{1}{s} \approx 10^{-5}\ \mathrm{GeV}^{-2} \approx 4000\ \mathrm{pb} \tag{6}$$

one obtains for a typical resonance of mass $M_R \approx 200$ GeV and width $\Gamma_{in} \approx 200$ MeV

$$\sigma_0 = 4\pi^2\,\frac{\Gamma_{in}}{M_R s} \approx 200\ \mathrm{pb}. \tag{7}$$

This cross section is three orders of magnitude larger than the detection limit of 0.1 pb, which corresponds to 10 events for an integrated luminosity of 100 pb^{-1} usually assumed to be achieved within one year.

In order to obtain the production cross sections for the various new particles shown in Fig. 1, one has to mulitply the cross section σ_0 with the corresponding parton densities. It is instructive to compare the different channels for a resonance with fixed mass and partial width to the initial two-parton state, which we choose to be $M_R = 200$ GeV $(x_R=0.4)$ and $\Gamma_{in} = 200$ MeV.

(i) Leptoquarks

As we will discuss in detail later, rare processes allow only leptoquark couplings to one electron (quark) chirality. Then the interactions (1a,b) ($e_{L,R} = \frac{1}{2}(1\mp\gamma_5)e$)

$$L_1 = \lambda_1\ \bar{u}^c e_{L,R}\ S + \text{h.c.},$$

$$L_2 = \lambda_2\ \bar{u} e_{L,R}\ R + \text{h.c.},$$

yield the cross sections

$$\sigma_S = \int_0^1 dx\ \sigma_S(\hat{s}) = \frac{1}{2}\ \sigma_o\ u(x_R)$$

$$\approx 50 \text{ pb}, \qquad (8)$$

$$\sigma_R = \int_0^1 dx\ \sigma_R(\hat{s}) = \frac{1}{2}\ \sigma_o\ \bar{u}(x_R)$$

$$\approx 10 \text{ pb}, \qquad (9)$$

where $u(x)$ ($\bar{u}(x)$) is the probability density of u-quarks (anti-u-quarks) in the proton which we have taken from ref. [6] with $Q^2 = 10^4$ GeV^2. Since protons contain more quarks than anti-quarks it is easier to produce leptoquarks with fermion number F=2 than those with fermion number F=0 ($F= \frac{1}{3}\cdot B+L$, where B and L are baryon and lepton number respectively). The converse is true for e^+p-collisions.

(ii) Excited leptons and quarks

Constraints from magnetic moments (cf. eq. (17)) allow also for excited fermions interactions with only one chirality of the ordinary "ground state" fermion, i.e. (cf. eqs. (1c,d)),

$$L_3 = \frac{\lambda_3}{\Lambda}\ \bar{u}^*\sigma^{\mu\nu}u_{L,R}\ F_{\mu\nu} + \text{h.c.},$$

$$L_4 = \frac{\lambda_4}{\Lambda}\ \bar{e}^*\sigma^{\mu\nu}e_{L,R}\ F_{\mu\nu} + \text{h.c.} \quad .$$

In order to obtain the u*- and e*- production cross sections the quark densities have to be convoluted with the probability distributions $\gamma(x)$ to obtain a photon with momentum fraction x from electron or quarks. Based on ref. [6] we estimate

$$\sigma_{u*} = \sigma_o \int_{x_R}^1 dxdx'\ \frac{1}{2}\ u(x')\gamma_e(x)\delta(xx'-x_R)$$

$$\approx 1 \text{ pb}, \qquad (10)$$

$$\sigma_{e*} = \sigma_o \sum_{f=q,\bar{q}} \int_{x_R}^1 dxdx'\ f(x)\gamma_f(x')\delta(xx'-x_R)$$

$$\approx 0.2 \text{ pb}. \qquad (11)$$

The u*- and e*- production cross sections are smaller than the leptoquark cross sections because they require a Bremsstrahlung photon. This is not the case for leptogluons which, if they should exist, will be copiously produced in ep-collisions due to the large gluon content of the proton. For the chiral interaction

$$L_5 = \frac{\lambda_5}{\Lambda}\ \bar{e}_8^A\sigma_{\mu\nu}e_{L,R}\ G_{\mu\nu}^A + \text{h.c.}$$

one obtains

$$\sigma_{e_8} = \sigma_o\ G(x_R) \approx 10 \text{ pb}, \qquad (12)$$

where $G(x)$ is the gluon distribution function (cf. ref. [6]).

(iii) $\gamma\gamma$- and γg-processes

The production of heavy scalar particles in $\gamma\gamma$- or γg-fusion is strongly suppressed because of the small parton luminosities. For the interactions L_6 and L_7 of eqs. (1f,g) a crude estimate (cf. ref. [6]) yields:

$$\sigma_{x_8} = \sigma_0 \int_{x_R}^{1} dx dx' \gamma_e(x) G(x') \delta(xx'-x_R)$$

$$< 0.1 \text{ pb}, \tag{13}$$

$$\sigma_\chi = \sigma_0 \sum_{f=q,\bar{q}} \int_{x_R}^{1} dx dx' \gamma_e(x) \gamma_f(x') \delta(xx'-x_R)$$

$$< 0.01 \text{ pb}. \tag{14}$$

From our discussion of the various two-parton channels we conclude that the only new heavy particles with sufficiently large production cross section are leptoquarks, and excited quarks and leptons, which we will study in more detail in the remaining part of this chapter.

Let us first turn to the production of excited quarks and leptons. As already mentioned, the conservation of the electromagnetic current forbids off-diagonal vector couplings, since otherwise one would have

$$\partial^\mu(\bar{e}^*\gamma_\mu e+\bar{e}\gamma_\mu e^*)$$

$$= -i(m_e-m_{e^*})\bar{e}^*e + \text{h.c.} + \ldots \neq 0. \tag{15}$$

Hence the allowed interaction to lowest order in $\frac{1}{\Lambda}$ is the gauge invariant, chirality changing tensor coupling (cf. (1d))

$$L_{e^*\gamma} = \frac{ef_\gamma}{2\Lambda} \bar{e}^*\sigma^{\mu\nu}(1+\chi\gamma_5)e F_{\mu\nu} + \text{h.c.}, \tag{16}$$

with $e = \sqrt{4\pi\alpha}$ and $f_\gamma = O(1)$. The coefficients in (16) are strongly constrained through the contribution of the excited electron to the electron anomalous magnetic moment (cf. Fig.3) which has to satisfy the bound $\delta(\frac{g_e-2}{2}) < 2\cdot 10^{-10}$. This implies [7]

$$1-|\chi|^2 < \frac{1}{f_\gamma^2} \frac{\Lambda}{22 \text{ TeV}} . \tag{17}$$

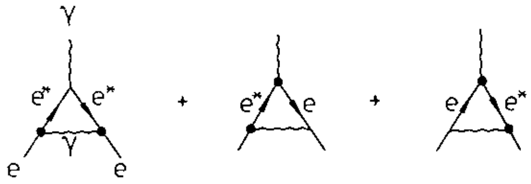

Fig. 3 Contribution of an excited electron (e*) to the electron anomalous magnetic moment

Furthermore the bound on the electric dipole moment of the electron, $d_e < 10^{-24}$ cm, yields the constraint [7]

$$\text{Im } \chi < \frac{1}{f_\gamma^2} \frac{\Lambda}{5\cdot 10^4 \text{ TeV}} . \tag{18}$$

With $f_\gamma = O(1)$ and $\Lambda = O(1 \text{ TeV})$ we conclude from (17) and (18) $\chi = \pm 1$, i.e., the excited electron e* can couple only to one chirality of the electron. If the interaction (16) involves mass eigenstates e and e*, there is clearly no constraint on the magnitude f in the case of a chiral coupling ($\chi = \pm 1$).

How natural is the required coupling of heavy lepton states to only one chirality of the light states? An excited electron with mass $m_{e^*} = O(\Lambda)$ has to form a real SU(2) x U(1) representation, such as

$$e_L^* \propto e_R^* \propto (1;-1) \quad \text{or} \tag{19a}$$

$$l_L^* \propto l_R^* \propto (2;-\tfrac{1}{2}),\ l^* = \begin{pmatrix} \nu^* \\ e^* \end{pmatrix} . \tag{19b}$$

Hence there exists one coupling to gauge fields of order $\frac{1}{\Lambda}$, whereas the second one is necessarily of order $\frac{1}{\Lambda^2}$. For instance, for the singlet case (19a) one has

$$L = \frac{g'f_L}{\Lambda} \bar{e}_L^* Y\sigma^{\mu\nu} e_R B_{\mu\nu}$$
$$+ \frac{g'f_R}{\Lambda^2} \bar{e}_R^* Y\sigma^{\mu\nu}\phi^+ l_L B_{\mu\nu} + \text{h.c.}, \tag{20}$$

where l_L is the ordinary lepton doublet, Y the U(1) hypercharge generator and $B_{\mu\nu}$ the corresponding field strength. After spontaneous SU(2) x U(1) breaking the second term in (20) gives a coupling to e_L of strength $\frac{g' f_R v}{\Lambda^2}$. Since it has the structure of an off-diagonal Yukawa coupling the phenomenologically required condition $f_R \frac{v}{\Lambda} \ll f_L$ appears indeed natural.

More problematic are excited leptons with masses $m_{e*} = O(100\ \mathrm{GeV}) \ll \Lambda$, which are relevant for HERA. In this case their masses should arise from spontaneous symmetry breaking and the "excited" leptons should form chiral representations of SU(2) x U(1). In the simplest case of "sequential leptons", i.e., $l_L^* \propto (2;-\frac{1}{2})$, $e_R^* \propto (1;-1)$, all tensor coupling have the structure of off-diagonal Yukawa couplings and are hence expected to be small. One could also consider the case of "mirror leptons", i.e., $l_L^* \propto (1;-1)$, $e_R^* \propto (2;-\frac{1}{2})$. In this case tensor couplings of order $\frac{1}{\Lambda}$ are possible, but their coefficients are likely to be small for the same reason which forbids mass terms of mirror leptons with ordinary leptons. We conclude that the phenomenologically required couplings of heavy lepton states to ordinary leptons of only one chirality are theoretically difficult to understand for excited lepton masses of order 100 GeV. The same applies to excited quarks.

Some preon models also predict the existence of leptogluons, i.e., colour-octet fermions which carry lepton number. Hence they interact with ordinary leptons and gluons:

$$L_{e_8 g} = \frac{g_s f_g}{2\Lambda} \bar{e}^{A}\sigma^{\mu\nu}(1\pm\gamma_5) e\, G^A_{\mu\nu} + h.c. \quad . \tag{21}$$

Like excited leptons leptogluons are allowed to interact with only one chirality of ordinary leptons.

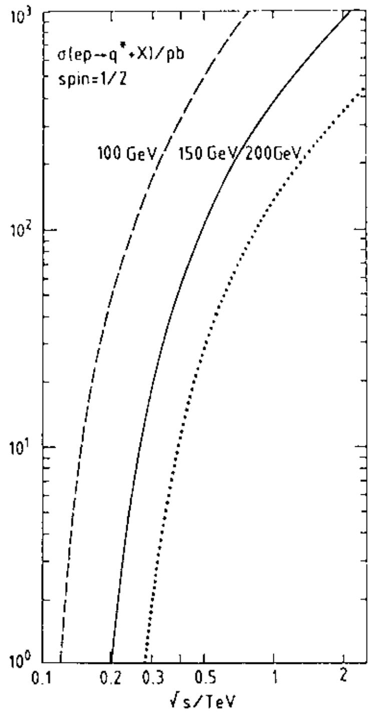

Fig. 4 q*-production cross section as function of $\sqrt{s}$ with the convention $\frac{f_\gamma}{\Lambda} = \frac{1}{m_{q*}}$ for three different masses m_{q*}. From ref. [9]

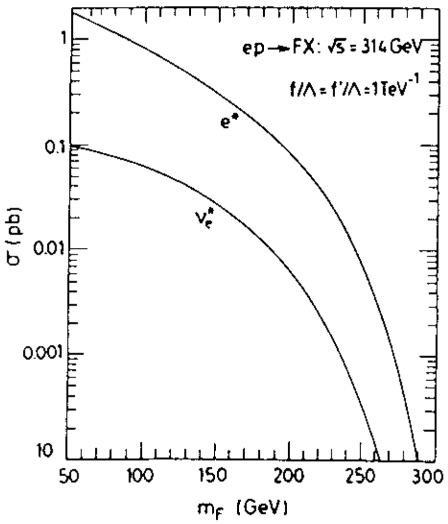

Fig. 5 e*-production cross section at $\sqrt{s}$=314 GeV and $\frac{f_\gamma}{\Lambda}$ = (1 TeV)$^{-1}$ as function of the excited lepton mass m_F. From ref. [10]

The production of excited fermions has been studied in considerable detail in the literature [1-3,8-13]. The partial widths of the heavy fermions to the corresponding initial states are:

$$\Gamma(F^* \to F\gamma) = \frac{\alpha}{4} \left(\frac{f_\gamma}{\Lambda}\right)^2 m^3_{F^*}, \quad (F^* = q^*, l^*), \qquad (22a)$$

$$\Gamma(e_8 \to eg) = \frac{\alpha_s}{4} \left(\frac{f_g}{\Lambda}\right)^2 m^3_{e_8}. \qquad (22b)$$

Fig. 4 shows the q*- production cross section as function of $\sqrt{s}$, as calculated by Kühn, Tholl and Zerwas [9], with the convention $\frac{f_\gamma}{\Lambda} = \frac{1}{m_{q^*}}$. For $\frac{f_\gamma}{\Lambda} = (1\ \text{TeV})^{-1}$ the detection limit of 0.1 pb is reached for $m_{q^*} \approx 250$ GeV. The e*-production cross section from the work of Hagiwara, Komamiya and Zeppenfeld [10] is shown in Fig. 5. About 50% of the cross section is due to the quasi-elastic process pe → pe*, which

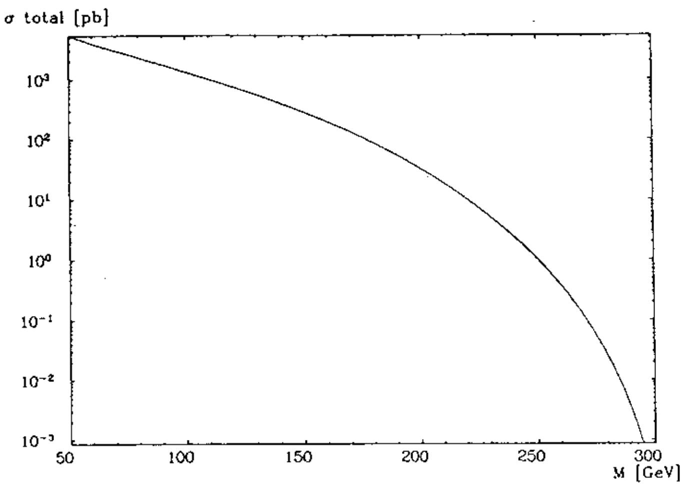

Fig. 6 Production cross section of leptogluons with Λ=1 TeV, f_g=1 (cf. eq. (21)). From ref. [12]

leads to the striking signal of a pure photon-electron final state, since the final state proton is not detected. The detection limit of 0.1 pb is reached for $m_{e^*} \approx 200$ GeV, in agreement with the results of Kühn et al. [9]. The e_8- production cross section of ref. [12] is shown in Fig. 6. It is larger than 0.1 pb for leptogluon masses below 250 GeV.

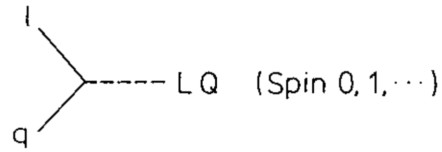

Fig. 7 Leptoquark coupling to lepton-quark pair

Leptoquarks are bosons which couple to quark-lepton pairs (cf. Fig. 7). Contrary to "excited" fermions leptoquarks of spin zero or one can have renormalizable couplings to quarks and leptons. Their existence is predicted by almost all extensions of the standard model, such as preon models or unified gauge theories. The most general SU(3) x SU(2) x U(1) invariant effective lagrangian [14] for leptoquarks with spin 0 or 1 allows the following SU(2) x U(1) quantum numbers:

$$S = 0:\ (2;\ \tfrac{1}{6}),\ (2;\ \tfrac{1}{7}),\ (3;\ \tfrac{7}{6}),\ (1;\ -\tfrac{1}{3}),\ (1;\ -\tfrac{4}{3}), \qquad (23a)$$

$$S = 1:\ (2;\ \tfrac{1}{6}),\ (2;\ -\tfrac{5}{6}),\ (3;\ \tfrac{2}{3}),\ (1;\ \tfrac{5}{3}),\ (1;\ \tfrac{2}{3}). \qquad (23b)$$

Of particular theoretical interest are scalar leptoquarks G^α with SU(2) x U(1) quantum number $(1;\ -\frac{1}{3})$. Together with Higgs doublets ϕ^i they form complete 5-plets of SU(5):

$$H^a = (G^\alpha,\ \phi^i),\ a = 1...5. \qquad (24)$$

In many models, in particular all supersymmetric extensions of the standard model, one has at least two Higgs multiplets H_1 and H_2 with the following couplings to the quark-lepton multiplets $\psi_a \propto \underline{5}^*$ and $\chi^{ab} \propto \underline{10}$:

$$L = \lambda_1 \psi_a^T C \chi^{ab} H^*_{1b} + \frac{1}{8}\lambda_2 \varepsilon_{abcde} \chi^{abT} C \chi^{cd} H_2^e + \text{h.c.} \quad . \qquad (25)$$

Using the decomposition

$$\psi_a = (d^c_{R\alpha}, \varepsilon_{ij} l_L^j), \qquad (26a)$$

$$\chi^{ab} = (\varepsilon^{\alpha\beta\gamma} u^c_{R\gamma}, q_L^{\alpha i}, \varepsilon^{ij} e_R^c), \qquad (26b)$$

one obtains in the low energy effective theory, where the SU(5) symmetry is broken to the subgroup SU(3)xSU(2)xU(1):

$$\begin{aligned} L = \lambda_1 (& d^{cT}_{R\alpha} C q_L^{\alpha i} \phi^*_{1i} - l_L^{iT} C e_R^{\,c} \phi^*_{1i} \\ & - \eta_1 \varepsilon^{\alpha\beta\gamma} d^{cT}_{R\alpha} C u^c_{R\beta} G^*_{1\gamma} - \eta_2 \varepsilon_{ij} l_L^{jT} C q_L^{\alpha i} G^*_{1\alpha}) \\ & + \lambda_2 (\varepsilon_{ij} u^{cT}_{R\alpha} C q_L^{\alpha i} \phi_2^j + \frac{1}{2}\eta_3 \varepsilon_{\alpha\beta\gamma} \varepsilon_{ij} q_L^{\alpha iT} C q_L^{\beta j} G_2^\gamma \\ & - \eta_4 u^{cT}_{R\alpha} C e^c_R G_2^\alpha) + \text{h. c.} \quad . \end{aligned} \qquad (27)$$

In ordinary unified theories with spontaneously broken SU(5) eq. (25) implies $\eta_1=\eta_2=\eta_3=\eta_4=1$. In this case the lagrangian (27) violates baryon number, and the observed stability of the proton requires very large leptoquark masses, $m_{LQ}>O(10^{10}\text{ GeV})$ [15]. However, there are examples of higher dimensional unified theories [16] where the Yukawa couplings in the effective low energy theory are not SU(5) symmetric. Leptoquark masses of order 100 GeV are phenomenologically only acceptable if baryon number (B) and lepton number (L) are conserved. Hence there exist two phenomenologically interesting cases:

$$\text{(i)} \quad \eta_1 = \eta_3 = 0, \quad \eta_2 = \eta_4 = 1,$$
$$B(G) = \frac{1}{3}, \qquad L(G) = 1 \; ; \qquad (28a)$$

$$\text{(ii)} \quad \eta_1 = \eta_3 = 1, \quad \eta_2 = \eta_4 = 0,$$
$$B(G) = -\frac{2}{3}, \qquad L(G) = 0. \qquad (28b)$$

In the first case the colour triplets $G^\alpha_{1,2}$ are leptoquarks which carry baryon and lepton number, in the second case they appear as diquarks with vanishing lepton number.

The possible couplings of SU(5)-type leptoquarks are further restricted by rare decays [17-19]. The strongest constraints exist for flavour changing couplings. For instance, the upper bound $BR(K^+ \to \pi^+ \text{ nothing}) < 1.4\cdot10^{-7}$ implies (cf. Fig. 8c)]

$$(\lambda_1)_{s_L\nu_e} (\lambda_1)_{d_L\nu_e} < \left(\frac{m_{G_1}}{10\text{ TeV}}\right)^2 . \qquad (29)$$

This bound can be satisfied by adjusting $(\lambda_1)_{s_L\nu_e}= 0$ or $(\lambda_1)_{d_L\nu_e}= 0$. However, SU(2) invariance then implies the existence of the flavour changing process $D^+ \to e^+e^-X$, where X denotes light hadrons. For instance, a bound for the branching ratio of 10^{-4} would yield the constraint

$$\frac{\sin\theta_C}{\cos\theta_C} (\lambda_1)^2_{u_L e_L} < \left(\frac{m_{G_1}}{1\text{ TeV}}\right)^2 \quad , \qquad (30)$$

where θ_C is the Cabibbo angle.

Interesting upper bounds have already been obtained for leptonic decay modes of D-mesons [20,21], and it is important to pursue the search for flavour changing decays of D- and B-mesons. Constraints on the mass mixing $m_{G_{1,2}}$ of the two leptoquarks G_1 and G_2 can be derived from leptonic decays of light mesons, since in general it leads to helicity unsuppressed decays (cf. Fig. 8b). For instance, from the process $\pi^+ \to e^+\nu_e$ one finds

$$(\lambda_1)_{u_L e_L} (\lambda_1)_{d_L\nu_e} \frac{m^2_{G_{1,2}}}{m_{G_1} m_{G_2}} < \frac{m_{G_1} m_{G_2}}{(10\text{ TeV})^2} \quad . \qquad (31)$$

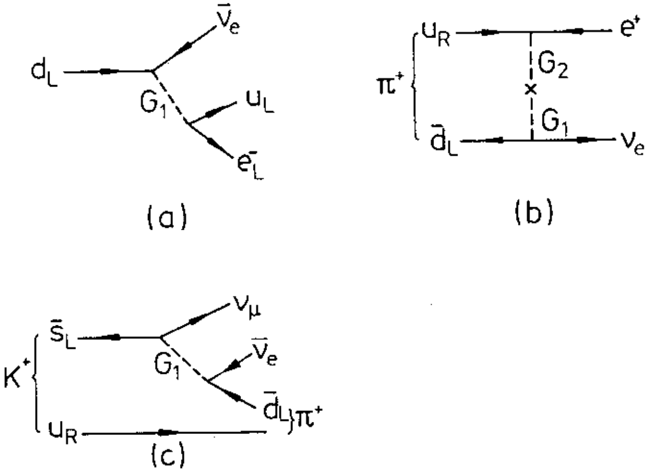

Fig. 8 Contribution from leptoquark exchange to low energy processes: (a) quark-lepton universality, (b) $\pi^+\to e^+\nu$, (c) $K^+\to\pi^+\nu\bar{\nu}$

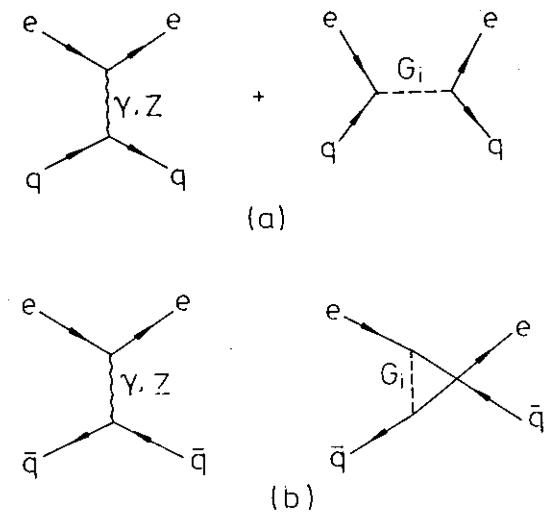

Fig. 9 s- and t-channel contributions of leptoquarks in the processes $eq\to eq$ and $e\bar{q}\to e\bar{q}$

Finally, leptoquarks lead to a deviation from quark-lepton universality in weak interactions, i.e., to a difference of the Fermi-constants measured in μ- and β-decays (cf. Fig. 8a). The experimentally observed equality of $G_F^{(\mu)}$ and $G_F^{(\beta)}$ yields the bound [17]

$$(\lambda_1)_{u_L e_L}\,(\lambda_1)_{d_L \nu_e} < \left(\frac{m_{G_1}}{2\ \mathrm{TeV}}\right)^2 . \qquad (32)$$

The constraints (29)-(32) on the leptoquark couplings are rather restrictive. However, it is conceivable that they are naturally satisfied in theories with a larger Higgs sector which contains two Higgs doublets and two leptoquarks for each quark-lepton family.

The production of leptoquarks in ep-collisions at HERA energies has recently been studied by a number of authors [14,22-25]. In the process $eq\to eq$ they contribute in the s-channel, in the process $e\bar{q}\to e\bar{q}$ in the t-channel (cf. Fig. 9). For the complete differential cross section one obtains [14]

$$\frac{d\sigma(e_{L,R}p)}{dx\,dQ^2} = \frac{1}{16\pi^2\hat{s}^2}\sum_q \Big\{ q(x,Q^2)\,|A(e_{L,R}q)|^2 + \bar{q}(x,Q^2)\,|A(e_{L,R}\bar{q})|^2 \Big\} , \qquad (33)$$

$$|A(e_{L,R}q)|^2 = |A_\gamma + A_Z|^2_{L,R} + 2\mathrm{Re}[(A_\gamma+A_Z)A_G^*]_{L,R} + |A_G|^2_{L,R} , \qquad (34a)$$

$$\begin{aligned} |A_\gamma + A_Z|^2_{L,R} = {} & \frac{2e^4 Q_e^2 Q_q^2}{t^2}(\hat{s}^2+u^2) \\ & + \frac{4e^2 g^2 Q_e Q_q}{t(t-m_Z^2)}(v_e \pm a_e)\,[v_q(\hat{s}^2+u^2) \pm a_q(\hat{s}^2-u^2)] \\ & + \frac{2g^4}{(t-m_Z^2)^2}(v_e\pm a_e)^2[(v_q^2+a_q^2)(\hat{s}^2+u^2) \\ & \pm 2v_q a_q(\hat{s}^2-u^2)] , \end{aligned} \qquad (34b)$$

$$\mathrm{Re}[(A_\gamma+A_Z)A_G^*]_{L,R} = -\frac{\lambda\hat{s}^2(\hat{s}-m_G^2)}{(\hat{s}-m_G^2)^2+m_G^2\Gamma_G^2}$$

$$\cdot\left[\frac{e^2Q_eQ_q}{t} + g^2\frac{(v_e\pm a_e)(v_q\pm a_q)}{t-m_Z^2}\right], \tag{34c}$$

$$|A_G|^2_{L,R} = \frac{1}{2}\frac{\lambda^2\ \hat{s}^2}{(\hat{s}-m_G^2)^2+m_G^2\Gamma_G^2}\ . \tag{34d}$$

Here $q(x,Q^2)$ is the probability density of finding a quark q with momentum fraction x inside the proton; $\hat{s}$,t and u are the usual Mandelstam variables of the parton subprocess, e and $g=e/(\sin\theta_W\cos\theta_W)$ are the gauge couplings of photon and Z-boson, Q_e and Q_q are the electric charges of electron and quark, and $v_{e,q}$ and $a_{e,q}$ denote the usual vector and axial vector couplings with the convention $a_e = -\frac{1}{4}$. The upper (lower) signs in the standard model amplitude (34b) and the interference term (34c) correspond to left-handed (right-handed) electrons and $\lambda=\lambda_1(\lambda_2)$, $G=G_1(G_2)$ in eqs. (34c,d). Expressions similar to eqs. (34) hold for the process $e_{L,R}\,\bar{q} \to e_{L,R}\bar{q}$.

The total cross section computed from (33) is shown as function of the leptoquark mass in Fig. (10) for the couplings $\lambda_1 = 0.3$ and $\lambda_1 = 0.1$, which are of the order of the bound (32) for $m_{G_1} = 0(200\ \mathrm{GeV})$. The corresponding differential

cross section, which is plotted in Fig. 11, shows a large, narrow peak at $x = \frac{m_G^2}{s}$. It is remarkable that for coupling strengths of the order of gauge couplings leptoquarks can be produced at HERA almost up to the phase-space boundary of 314 GeV. For smaller masses they can be discovered for much smaller couplings, e.g. $\frac{\lambda_1^2}{4\pi} = 10^{-3}\ \alpha_{em}$ for m_{LQ}=200 GeV [14].

In addition to the simplest process ep → GX, Dobado, Herrero and Muñoz [24] have considered the processes ep → GgX, ep → GuX and ep → GG*X. Also interesting are "t-channel" leptoquarks which could be produced together with t-quarks in electron-gluon fusion. Their characteristic features have been discussed elsewhere [1,3,26,27].

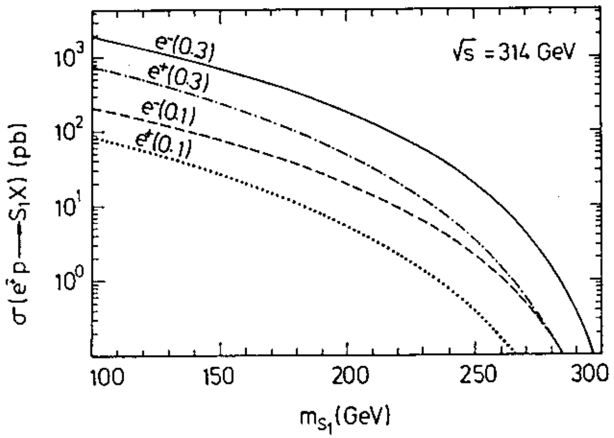

Fig. 10 Total production cross section of leptoquarks at $\sqrt{s}$ = 314 GeV in e^-p- and e^+p-collisions (S_1 = G, λ_1 = 0.1, 0.3; cf. eq. (34)). From Ref. [14].

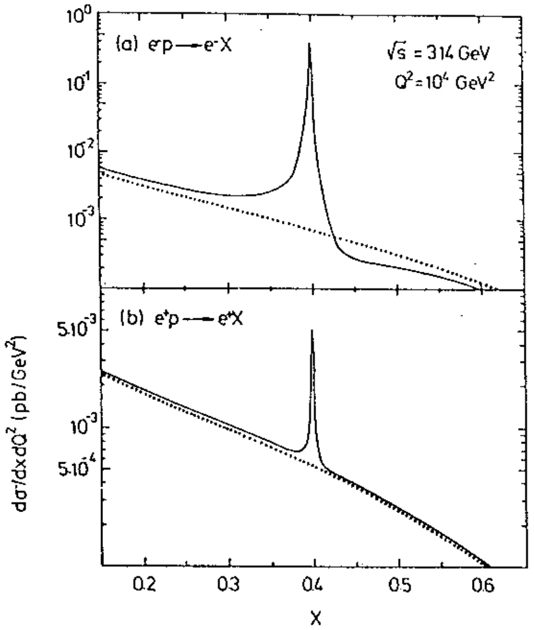

Fig. 11 Differential cross section for leptoquark production at $\sqrt{s}$ = 314 GeV and $Q^2=10^4$ GeV^4 with m_G=200 GeV and λ_1= 0.3. From ref. [14].

3) Pair production of new particles

New particles with a new, conserved quantum number can only be produced in pairs. This is the case for almost all supersymmetric extensions of the standard model on which we will concentrate in this section. The new s-particles are charged scalars for each chirality of lepton and quark fields, gaugions, and additional neutral and charged higgsinos, which

are present because the minimal supersymmetric standard model requires two Higgs-doublet superfields H_1 and H_2. Our notation reads:

$$
\begin{array}{lcl}
l = l_L + l_R & : & \tilde{l}_L,\ \tilde{l}_R \\
q = q_L + q_R & : & \tilde{q}_L,\ \tilde{q}_R \\
\gamma, Z & : & \tilde{\gamma}, \tilde{Z}, \tilde{h}_1, \tilde{h}_2 \\
W^+ & : & \tilde{w}^+,\ \tilde{h}^+ \\
g & : & \tilde{g} \quad .
\end{array}
$$

The new, conserved quantum number is the discrete R-parity, which is zero for ordinary particles and nonzero for the new s-particles.

A specialfeature of supersymmetric theories is the appearance of Majorana fermions, satisfying the condition

$$\psi = \psi^C = C\bar{\psi}^T \ , \tag{35}$$

where C is the charge conjugation matrix. The four "neutralinos" $\tilde{\gamma}, \tilde{z}, \tilde{h}_1$ and $\tilde{h}_2$ are such Majorana fermions. Since they carry no quantum numbers they can have a general mass matrix

$$
\begin{aligned}
L_M &= -\frac{1}{2}\,\mu_{ij}\,\bar{\lambda}_i\lambda_j \\
&= -\frac{1}{2}\,\mu_{ij}\,\lambda_i^T C\lambda_j \ .
\end{aligned} \tag{36}
$$

Their interactions with ordinary fermions can be written in the usual form (cf. Fig. 12a)

$$L \ = \ -ig\ \tilde{f}^*_{L,R}\ \bar{\lambda} f_{L,R} + \text{h.c.}, \tag{37}$$

where $\tilde{f}_{L,R}$ is the complex scalar field associated with the fermion $f_{L,R}$. However, because of (35) three propagators

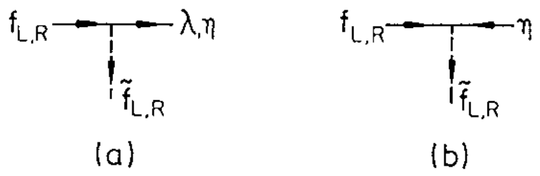

Fig. 12 Couplings of "neutralinos" (λ) and "charginos" (η) to ordinary left- and right-handed fermions ($f_{L,R}$) and the associated charged scalar particles ($\tilde{f}_{L,R}$)

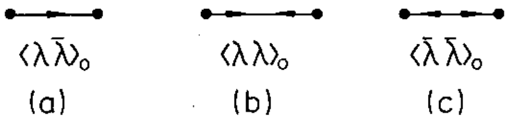

Fig. 13 Propagators of Majorana fermions

appear in the Feynman rules, $\langle\lambda\bar{\lambda}\rangle_0$, $\langle\lambda\lambda\rangle_0$ and $\langle\bar{\lambda}\bar{\lambda}\rangle_0$ (cf. Fig. 13). For charged gauginos η, which do not satisfy eq. (35), only the first propagator appears. They have, however, two kinds of vertices with ordinary fermions $f_{L,R}$, $f'_{L,R}$ (cf. 12 a,b)

$$L_I^{(1)} = -ig\ \tilde{f}^*_{L,R}\ \bar{\eta} f_{L,R} + \text{h.c.}\ , \tag{38a}$$

$$L_I^{(2)} = -ig'\ \tilde{f}'^*_{L,R}\ \bar{\eta}^c f'_{L,R} + \text{h.c.} \tag{38b}$$

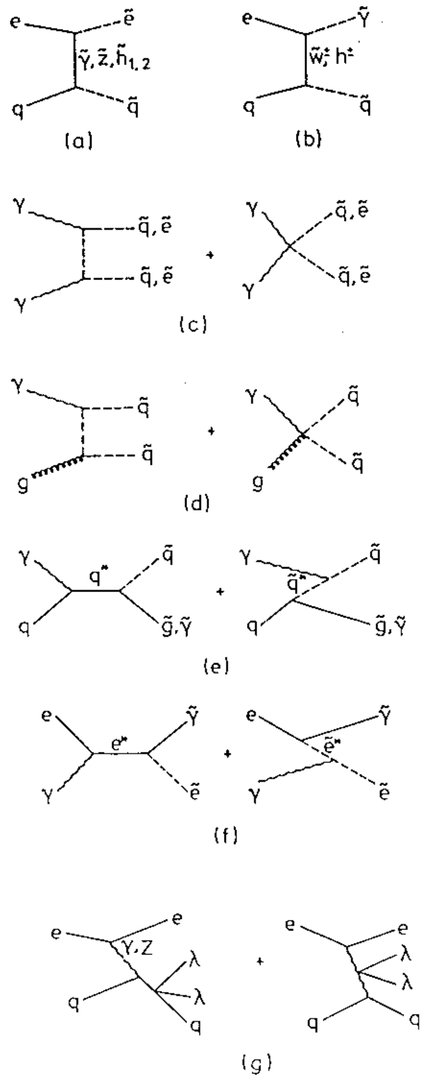

Fig. 14 Pair production of supersymmetric particles in different two-parton processes in electron-proton scattering.

Given the Feynman rules one can easily write down the various processes for the production of supersymmetric particles in inelastic electron-proton scattering (cf. Fig. 14). The most important one is the pair production of scalar quarks and leptons (Fig. 14a-d), then there are the Compton-type processes for the production of gaugino- scalar quark (lepton) pairs (Fig. 14e,f) and finally the pair production of goldstinos (Fig. 14g).

Of particular interest is the pair production of scalar quarks and leptons through neutral gaugino and higgsino exchange (cf. Fig. 15) which has first been studied by Jones and Llewellyn Smith [28]. Since the scalar quark and lepton masses $m_{\tilde{l}}$ and $m_{\tilde{q}}$ cannot be neglected the range in x and Q^2 is kinematically restricted:

$$x_- \leq x \leq 1, \tag{39a}$$

$$x_- = \frac{1}{2}(m_{\tilde{l}}+m_{\tilde{q}})^2 , \tag{39b}$$

$$Q_-^2 \leq Q^2 \leq Q_+^2 , \tag{39c}$$

$$Q_\pm^2 = \frac{1}{2}\Big[\hat{s}-m_{\tilde{l}}^2-m_{\tilde{q}}^2 \pm \left((\hat{s}-m_{\tilde{l}}^2-m_{\tilde{q}}^2)^2 - 4m_{\tilde{l}}^2 m_{\tilde{q}}^2\right)^{\frac{1}{2}}\Big], \quad \hat{s} = xs. \tag{39d}$$

The total production cross section is obtained in the usual way from the differential parton cross section [28]

$$\sigma = \int_{x_-}^{1} dx \int_{Q_-^2}^{Q_+^2} dQ^2 \; \frac{1}{4} \sum_q \frac{d\hat{\sigma}}{dxdQ^2}\, q(x,Q^2), \tag{40a}$$

where

$$\frac{d\hat{\sigma}}{dxdQ^2} = \frac{1}{16\pi\hat{s}^2}\begin{cases} \left| \sum_i \frac{g_{l_i} g_{q_i} \mu_i}{Q^2 + \mu_i^2} \right|^2 & (LL),(RR) \\ \frac{1}{s}(Q_+^2-Q^2)(Q^2-Q_-^2)\left|\sum_i \frac{g_{l_i} g_{q_i}}{Q^2+\mu_i^2}\right|^2 & (LR),(RR) \end{cases} \tag{40b}$$

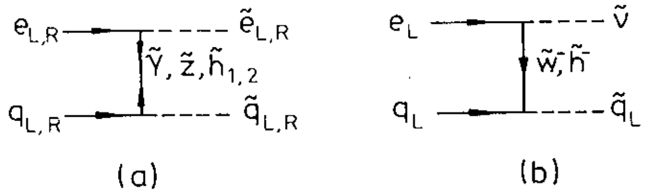

Fig. 15 Production of scalar quarks and leptons through "neutralino"- and "chargino"-exchange

Here μ_i, g_{l_i} and g_{q_i} are "neutralino" masses and couplings to leptons and quarks. (LL), (LR), (RL) and (RR) denote the four possible helicity configurations of the initial electron-quark state. The propagator $\langle\bar{\lambda}\bar{\lambda}\rangle_0$ of the exchanged neutralino differs from $\langle\lambda\bar{\lambda}\rangle_0$ by the charge conjugation matrix which commutes with the chirality projection operators. Hence the (LL) and (RR) amplitudes are proportional to the gaugino masses μ_i.

The size of the production cross section depends strongly on the neutralino mass matrix [29-32] which in the most general supersymmetric extension of the standard model takes the form:

$$M = \begin{pmatrix} m_{\tilde{\gamma}} & m_{\tilde{\gamma}\tilde{z}} & 0 & 0 \\ m_{\tilde{\gamma}\tilde{z}} & m_{\tilde{z}} & im_Z & 0 \\ 0 & im_Z & -\mu\sin2\theta_v & \mu\cos2\theta_v \\ 0 & 0 & \mu\cos2\theta_v & \mu\sin2\theta_v \end{pmatrix} . \tag{41}$$

Here $\tan\theta_v = \frac{v_2}{v_1}$, where $v_{1,2}$ are the vacuum expectation values of the two Higgs doublets; μ is the Dirac mass of the two higgsinos, and the $\tilde{\gamma}$-$\tilde{z}$ submatrix is determined by the two $U(1)_Y$ and $SU(2)_L$ gaugino masses m_1 and m_2. Under the simplifying assumptions

$$v_1 = v_2, \tag{42a}$$

$$\frac{3}{5}\cos^2\theta_W\, m_1 = \sin^2\theta_W m_2, \tag{42b}$$

two parameters remain which can be chosen to be the lightest "neutralino" (λ_1) and "chargino" (η_1) masses [30].

The results of Komatsu and Rückl [30] for the production cross sections $e^{\pm}p \rightarrow \tilde{e}\tilde{q}X$ and $e^{\pm}p \rightarrow \tilde{\nu}\,\tilde{q}X$ are shown in Fig. 16 as functions of $m_{\tilde{\ell}}+m_{\tilde{q}}$. The curves a to c represent different gaugino mixings, which can change the cross section up to a factor of 5. Case c, for instance, corresponds to $\lambda_1 \approx \gamma$, m_{λ_1}= 20 GeV, m_{η_1}=50 GeV; case b corresponds to $\lambda_1 \approx h_2$, m_{λ_1}= 20 GeV, m_{η_1}=30 GeV. According to Fig. 16 the detection limit of 0.1 pb is reached at HERA for scalar quark and lepton masses $m_{\tilde{\ell}}+m_{\tilde{q}} \approx$ 180 GeV.

In the minimal supergravity model the various supersymmetry breaking masses are all determined by two parameters, a universal gaugino mass $m_{\frac{1}{2}}$ and the gravitino mass m_o. Fig. 17, taken from ref. [30], shows that the range of parameters, which will be probed at HERA is a significant extension of present bounds on supersymmetric models.

Supersymmetric neutral current events can be distinguished from ordinary ones through missing energy and momentum (cf. Fig. 18), which is assumed to be carried away by photinos appearing in scalar quark and lepton decays: $\tilde{q}\rightarrow q\gamma$, $\tilde{e}\rightarrow e\gamma$. Monte Carlo studies then show that supersymmetric events can be clearly distinguished from ordinary events in the $\Delta\phi$-Δy plane [1]. Here $\Delta\phi$ is the acoplanarity angle (cf. Fig. 18) and

$$\Delta y = y_q - y_e,$$

$$y_q = \sum_i \frac{(E_i - p_{\parallel i})}{2E_e}\,, \quad y_e = \frac{E_e' - p_{\parallel e}'}{2E_e}\,, \qquad (43)$$

i.e., Δy is the difference between the usual variable $y = \frac{Q^2}{xs}$ determined from current jet and scattered electron. A detailed study of the identification of supersymmetric events has recently been performed by Stanco [33].

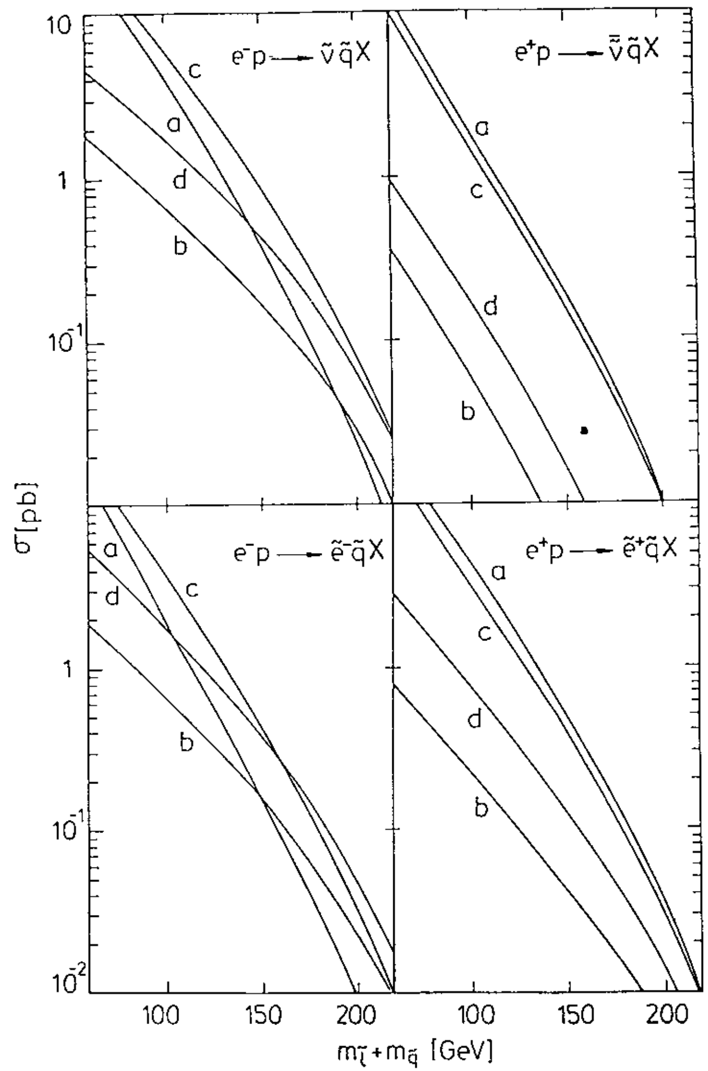

Fig. 16 Cross sections for pair production of scalar quarks and leptons as function of $m_{\tilde{\ell}}+m_{\tilde{q}}$ at $\sqrt{s}$ = 314 GeV ($m_{\tilde{\ell}_L}=m_{\tilde{\ell}_R}$, $m_{\tilde{q}_L}=m_{\tilde{q}_R}$). The curves a to c correspond to different "neutralino" and "chargino" mixings (see text). From ref [30]

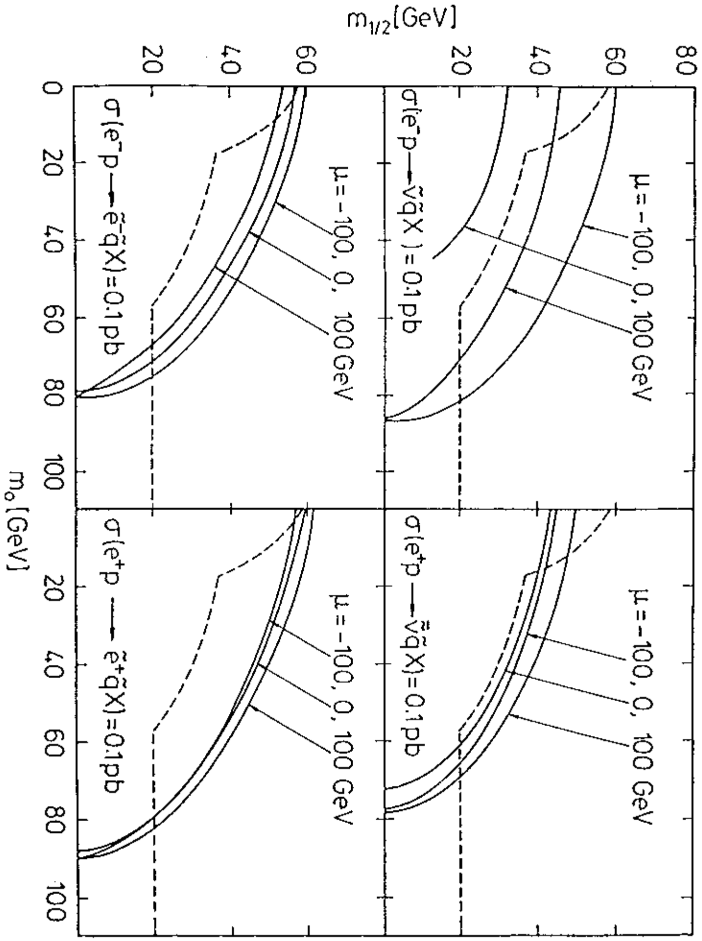

Fig. 17 Range of parameters of minimal supergravity model, which can be probed at HERA. The regions below the dashed curves are already excluded by present bounds from e^+e^- - and $p\bar{p}$ -colliders.

The supersymmetric Compton-type processes (cf. Fig. 14e,f) are interesting in the case of light photinos. Unfortunately the cross section is suppressed by one power of α_{em} compared to pair production of scalar quarks and leptons. The explicit calculation [34,35] shows indeed that for scalar electron and photino masses, which are not yet excluded by e^+e^- experiments, the production cross section is smaller than the detection limit of 0.1 pb.

Let us finally discuss pair production of goldstinos at HERA. Theories with spontaneously broken global supersymmetry contain a massless Goldstone fermion, the goldstino. At energies small compared to the masses of the heavy s-particles the interactions of goldstinos with ordinary matter are described by an effective nonrenormalizable lagrangian (cf. Fig. 19) which is invariant under nonlinear supersymmetry transformations. This invariance determines entirely the goldstino interactions at low energies [36]:

$$L_{eff} = L_0 + \frac{1}{2}\,\bar{\lambda}(i\not\partial - m_\lambda)\lambda + i\kappa\bar{\lambda}\gamma^\nu\partial^\mu\lambda\; T_{\mu\nu} + \ldots, \tag{44}$$

where L_0 is the lagrangian of ordinary matter, $T_{\mu\nu}$ the corresponding energy momentum tensor and m_λ an explicit supersymmetry breaking goldstino mass respectively. The parameter κ with mass dimension -4 is the analogue of f_π in the familiar chiral-invariant σ-models.

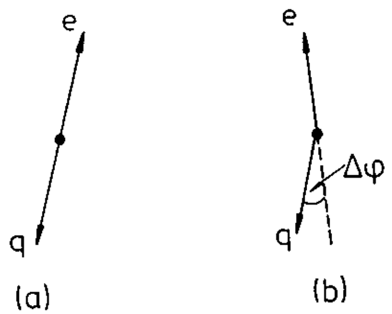

Fig. 18 Electron-quark configurations in the transverse plane from ordinary (a) and supersymmetric (b) neutral current processes

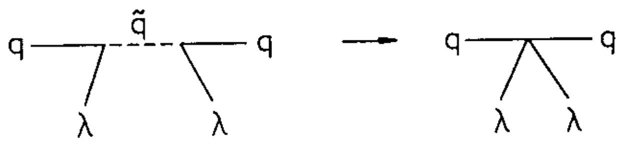

Fig. 19 Effective, nonrenormalizable quark-goldstino interaction at energies small compared to the scalar quark mass $m_{\tilde{q}}$

In a recent paper Ma, Nachtmann and Schücker have studied goldstino pair production at HERA [37]. The parton subprocess is

$$eq \to eq\ \lambda\lambda, \tag{45}$$

where the two goldstinos can be produced from every line of the basic scattering process eq→eq since they couple to the energy momentum tensor (cf. Fig. 14g). The corresponding observed process is (cf. Fig. 20)

$$\begin{aligned} e(E_e) + p(E_p) \to\ & e(E_e', \theta_e) + \text{quark jet } (E_T, \theta_q) \\ & + \text{spectator jet} + \lambda + \lambda\ . \end{aligned} \tag{46}$$

The two goldstinos escape undetected, and in order to distinguish the process (46) from a charged current event, where a neutrino escapes, it is important to detect the scattered electron. In order to simplify the calculation Ma et al. have restricted their analysis to events with small

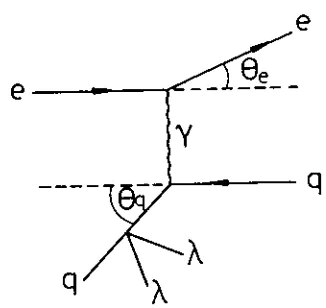

Fig. 20 Goldstino-pair production in electron-quark scattering

electron scattering angle ($\theta_e^{max} = O(1$ mrad$)$), i.e., to the Compton-type process $\gamma q \to \gamma q\ \lambda\lambda$. The predicted number of events is plotted in Fig. 21 as function of x and m_λ for an integrated luminosity of 200 pb^{-1}. For m_λ < 50 GeV values of $(x)^{-\frac{1}{4}}$ up to 50 GeV can be explored at HERA. Events with two goldstinos in the final state are very similar to events with a pair of scalars $\tilde{e}$ and $\tilde{q}$ which decay into $e\tilde{\gamma}$ and $q\tilde{\gamma}$. However, it should be possible to distinguish the two processes by means of their different differential cross sections.

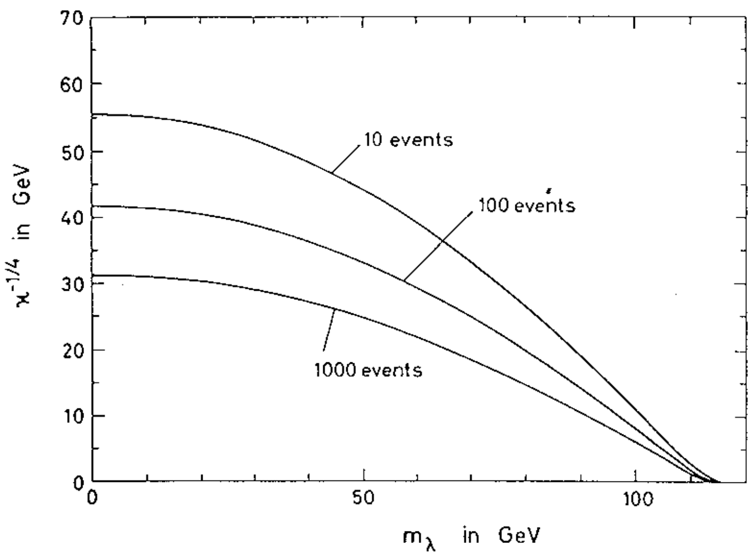

Fig. 21 Event rates for goldstino pair production with quark jet parameters E_T>2 GeV,10^o<θ_q<160^o and $\int$Ldt = 200 pb^{-1}. From ref. [37]

4. Helicity changing contact interactions

The standard model of strong and electroweak interactions is known to describe correctly the structure of matter for distances larger than the Fermi scale, i.e., $r > G_F^{\frac{1}{2}} \approx 10^{-16}$ cm. Low energy effects of new interactions, characterized by a mass scale $\Lambda > 1/\sqrt{G_F}$, can be systematically studied by means of an effective, nonrenormalizable lagrangian

$$L_{eff} = L_o + \frac{1}{\Lambda} L_1 + \frac{1}{\Lambda^2} L_2 + \ldots , \tag{47a}$$

where

$$L_o = L_{gauge} + L_{kin}$$

$$+ (\Gamma_u \bar{q}_L \phi u_R + \Gamma_d \bar{q}_L \phi d_R + \Gamma_e \bar{l}_R \phi e_R + h.c.). \tag{47b}$$

Here L_o is the standard model lagrangian, q_L, u_R, d_R, l_L, e_R and ϕ are quark-, lepton- and Higgs-fields respectively, and $\Gamma_{u,d,e}$ are the corresponding Yukawa couplings. L_n is the general linear combination of all SU(3)xSU(2)xU(1) invariant operators of dimension (4+n) which one can construct from the fields contained in L_o. If one imposes baryon- and lepton-number conservation, which is necessary for interaction scales Λ of order 1 TeV, the expansion (47) begins with dimension-6 operators. From a general analysis of rare processes [19,38] one obtains the bound $\Lambda > O(1$ TeV$)$ for flavour conserving operators, and $\Lambda > O(100$ TeV$)$ for flavour changing operators.

Flavour conserving four-fermion electron-quark operators can be probed in ep-collisions at HERA. One usually studies the effects of the various possible current-current operators [39]:

$$L_{eff} = \frac{g^2}{\Lambda^2} [\eta_{LL} \bar{e}_L \gamma^\mu e_L \bar{q}_L \gamma_\mu q_L + \eta_{LR} \bar{e}_L \gamma^\mu e_L \bar{q}_R \gamma_\mu q_R$$

$$+ \eta_{RL} \bar{e}_R \gamma^\mu e_R \bar{q}_L \gamma_\mu q_L + \eta_{RR} \bar{e}_R \gamma^\mu e_R \bar{q}_R \gamma_\mu q_R], \tag{48}$$

$$g^2 = 4\pi , \quad q = (u,d).$$

The scattering amplitude obtained from this lagrangian interferes with the standard model amplitude from γ- and Z-exchange, and a quantitative analysis shows that the effects of the contact interactions (48) will be observable at HERA for values of Λ up to (5-10) TeV [39,40].

The Ansatz (48), which conserves chirality at lepton and quark vertex, is natural since the smallness of the lepton and quark masses compared to Λ must be due to some custodial chiral symmetry. However, in the standard model this is the local SU(2)xU(1) invariance which allows also operators which flip the chiralities at lepton and quark vertex. The most general SU(3)xSU(2)xU(1) invariant, chirality changing four-fermion electron-quark operator [41] is given by

$$L_{eff}^{HF} = \frac{g^2}{\Lambda^2} \left[c_1(\bar{l}_L e_R)(\bar{d}_R q_L) + c_2(\bar{l}_L e_R)(\bar{q}_L u_R)\right.$$

$$\left. + c_3(\bar{l}_L u_R)(\bar{q}_L e_R) + h.c.\right]. \qquad (49)$$

For parameter values $\Lambda = O(1 \text{ TeV})$, $c_i = O(1)$, constraints from rare decays require [41] $c_1 = 0$, $c_2 = \frac{c_3}{2}$, for which (49) becomes a pure tensor interaction.

The chirality changing, SU(2)xU(1) invariant lagrangian (49) is interesting, because it gives rise to a qualitatively new effect. Contrary to all current-current interactions it generally leads to a violation of the Callan-Gross relation in electron-proton scattering. The structure functions $F_i(x,Q^2)$ ($i=1\ldots3$) for inelastic electron-proton scattering are defined through

$$\frac{d\sigma}{dxdy} = \frac{4\pi\alpha^2}{x^2y^2s}\left[y^2xF_1(x,Q^2)+(1-y)F_2(x,Q^2)\pm(y-\frac{y^2}{2})xF_3(x,Q^2)\right], \qquad (50a)$$

where

$$s = (p_e+p_p)^2, \quad -(p_e'-p_e)^2 = Q^2 = xys, \quad y = \frac{(p_e'-p_e)p_p}{p_e p_p}.$$

One easily verifies that for arbitrary current-current interactions the structure functions F_1 and F_2 satisfy the Callan-Gross relation [42]

$$R = \frac{2xF_1-F_2}{2xF_1} = 0. \qquad (51)$$

Scalar partons inside the proton would contribute only to F_2 and not to xF_1 and hence give a negative contribution to R, i.e., $\Delta R_S < 0$. On the contrary the contact interactions (49) give a positive contribution to R ($c_1=0$, $c_2 = \frac{c_3}{2} = 1$):

$$R(x,Q^2) = \frac{\frac{3}{2\alpha^2}\frac{Q^4}{\Lambda^4}\, x\,(u(x,Q^2)+\bar{u}(x,Q^2))}{F_2^{SM}(x,Q^2)+\frac{5}{2\alpha^2}\frac{Q^4}{\Lambda^4}\,x(u(x,Q^2)+\bar{u}(x,Q^2))}, \qquad (52)$$

where $F_2^{SM}(x,Q^2)$ and $u(x,Q^2)$ are the standard model structure function and the u-quark distribution in the proton respectively. Approximating F_2^{SM} by the electromagnetic part with $u(x) \approx 2d(x)$ the ratio R becomes independent of x:

$$R(Q^2) = \frac{\frac{3}{\alpha^2}\frac{Q^4}{\Lambda^4}}{1+\frac{5}{\alpha^2}\frac{Q^4}{\Lambda^4}}. \qquad (53)$$

This is also approximately the case for the full expression (52) which is plotted in Fig. 22 for $Q^2=10^4$ GeV^2 and two values of Λ. On the contrary QCD corrections yield significant deviations from R=0 only at small values of x [43]. Fig. 23 shows $R(x,Q^2)$ as function of Q^2 at x=0.5. An error on R of 0.1 at $Q^2=5\cdot10^3$ GeV^2 would correspond to a bound on the new interaction scale $\Lambda > 2$ TeV [44].

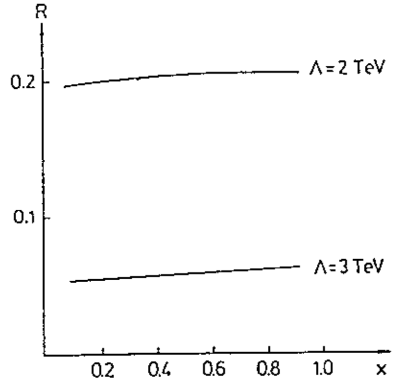

Fig. 22 x–dependence of R at $Q^2 = 10^4$ GeV² for two values of Λ

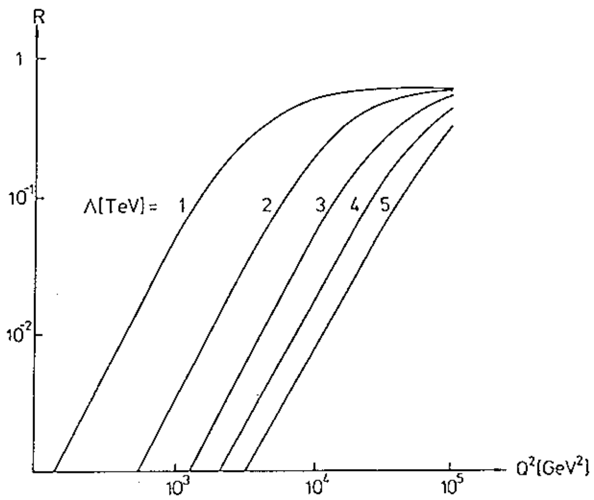

Fig. 23 Q²-dependence of R at x = 0.5

We conclude that an accurate determination of structure functions at HERA will provide non-trivial bounds on chirality changing contact interactions. However, from a theoretical point of view, the pure tensor structure of the chirality-flip contact terms, required by consistency with rare decays, is difficult to understand, and the observation of a deviation from the Callan-Gross relation at HERA energies would be surprizing.

5. Summary

In the previous sections we have surveyed some possibilities to search for new particles and to test for new interactions in electron-proton collisions at HERA. Among the various particles, which can be produced as s-channel resonances in a two-parton subprocess, leptoquarks play a special role, since they are the only new particles with renormalizable couplings to the available partons. Their production cross section is large. For coupling strengths of the order of gauge couplings leptoquarks can be discovered at HERA almost up to the phase-space boundary, i.e.,

$$m_{LQ} < 300 \text{ GeV}.$$

For smaller masses much smaller couplings can be detected, e.g., $\frac{\lambda^2}{4\pi} = 10^{-3}\,\alpha_{em}$ for $m_{LQ} = 200$ GeV. Leptoquarks would appear as narrow peaks in the x-distribution of the inelastic differential electron-proton cross section. Leptogluons, which can be formed in electron-gluon fusion, have a similar signature. Contrary to leptoquarks the production cross section is the same in e^-p- and e^+p- scattering. At HERA leptogluons with masses

$$m_{e_8} < 250 \text{ GeV}$$

could be detected. In contrast to leptoquarks the discovery of leptogluons at HERA would imply the existence of quark-lepton substructure with an interaction scale $\Lambda = O(1\ \text{TeV})$. The same is true for excited electrons and quarks which could be discovered at HERA in the mass range

$$m_{e*} < 200\ \text{GeV},$$

$$m_{q*} < 250\ \text{GeV}.$$

About 50% of the e*-production cross section is due to the quasi-elastic process $ep \to e^*p$ which leads to the very clean observed final state of an electron-photon pair with high transverse momentum.

The dominant process for pair production of supersymmetric particles is the exchange of neutral and charged gauginos and higgsinos. Pairs of scalar quarks and electrons can be detected in the mass range

$$m_{\tilde{e}} + m_{\tilde{q}} < 180\ \text{GeV}.$$

In the minimal supergravity model this corresponds to a range of parameters $m_{\frac{1}{2}} < 60$ GeV, $m_0 < 80$ GeV, where $m_{\frac{1}{2}}$ and m_0 denote a universal gaugino mass and the gravitino mass respectively. Models with nonlinearly realized global supersymmetry predict goldstino-pair production at HERA. This can be detected for goldstino "decay-constants" $\chi^{-\frac{1}{4}} < 50$ GeV.

Effects of new interactions beyond the standard model can be studied in a model-independent way by means of an effective lagrangian analysis. Four-fermion electron-quark operators of current-current type conserve chirality at electron and quark vertex. The corresponding scattering amplitudes interfere with the standard model γ- and Z-exchange amplitudes, and at HERA bounds on new interaction scales

$$\Lambda = (5\text{-}10)\ \text{TeV}$$

can be obtained. For chirality-flip, $SU(3)\times SU(2)\times U(1)$ invariant interactions bounds on Λ of about 2 TeV appear feasible.

The examples of new particles and interactions, which we have discussed in these lectures, show to what extent expectations for "new physics" are already limited by the success of the standard model. They also illustrate, however, that the exploration of the next order of magnitude in momentum transfer in electron-proton collisions could well lead to discoveries not yet anticipated by theorists.

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