preprint

Daresbury Laboratory

DL/SCI/P464A

BEAM LIFETIME ON THE DARESBURY SRS

by

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Presented to the 1985 Particle Accelerator Conf. Vancouver B.C., Canada,
13-16 May, 1985

MAY, 1985

Science and Engineering Research Council

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The derivation of what is thought to be a novel expression for beam decay based on a model including gas desorption is given and its predictions compared to observed beam decays in the SRS. The model is shown to predict beam decay to within the accuracy of observation which represents a significant improvement over the simple exponential decay model. The SRS beam decay has been observed for a period of many months and statistical correlation used to separate variations in zero-beam lifetime (due to vacuum base pressure) from the faster decay due to radiation desorbed gas. Evidence is presented to show that lifetimes which do not correlate are indicative of machine malfunction such as low RF power or misaligned closed orbit. Therefore this decay analysis is an extremely useful indicator of the "health" of the storage ring. The recovery of the unbaked machine from a complete vacuum let-up has also been recorded in detail and this history is illustrated.

Introduction

As in all Storage Rings, the decay of the stored beam in the Daresbury SRS is a subject of intense interest. A model for this decay process has been developed at Daresbury which predicts beam decay to an accuracy within the limits of experimental error. The parameters associated with the model have given some fascinating insights into machine operation.

Beam Decay Model

The presumption is made that if the RF over-voltage and closed orbit are kept constant then any variation in the rate of decay of a beam is due solely to a change in vacuum pressure. The decay can then be modelled by:

$$\frac{dI_b}{dt} = -K I_b P \quad (1)$$

where $\frac{dI_b}{dt}$ is the rate of beam decay at current I_b
P is the vacuum pressure at current I_b
K is a constant

From beam-induced gas desorption mechanisms one could speculate that

$$P = P_0 + A I_b \quad (2)$$

when P_0 is P at $I_b = 0$
A is a constant

Then

$$\frac{dI_b}{dt} = -K I_b (P_0 + A I_b) \quad (3)$$

which can be re-written

$$\frac{dI_b}{dt} = -I_b \left(\frac{1}{T_0} + B I_b \right) \quad (4)$$

where

$$T_0 = \frac{1}{K P_0} = \text{exponential decay constant at } I_b = 0$$

$$B = \frac{A}{T_0 P_0}$$

This may be integrated to give

$$I_b = \frac{I_0 e^{-t/T_0}}{1 + B T_0 I_0 (1 - e^{-t/T_0})} \quad (5)$$

which is the primary model for beam decay.

A typical fit of this model to an actual beam is shown in Figure 1 together with a simple exponential fit for comparison.

Figure 1. Actual and predicted beam current

The real and model curves are so close as to be indistinguishable in this figure size but the difference between the curves is plotted in Figure 2.

The standard deviation of the residual errors for the model is $\sim 0.3\%$ of the mean beam current compared with $\sim 2\%$ for the simple exponential curve.

Obviously this is a very good model. However it can be improved still further from examination of the dynamic behaviour of the storage ring vacuum pressure. This shows a rapid response to a change in beam conditions followed by a much slower exponential response. These presumably correspond to two different outgassing mechanisms, namely radiation desorption and thermal desorption. As a model for this process the following phase retard Laplace Transform was used.

$$P(s) = A k \left[\frac{1 + s T_v}{k + s T_v} \right] I_b(s) \quad (6)$$

Where:

k is the photo desorption fraction of a pressure change which must be greater than zero
A is the outgassing constant
 T_v is the vacuum chamber time constant
s is the Laplace operator.

Figure 2. Difference between actual beam current and the model and simple predictions

If, for simplicity, $I_b(s)$ is written as a simple exponential decay then the following expression for the time-varying vacuum pressure $P(t)$ is obtained:

$$P(t) = P_0 + AkI_b(t) + \frac{AT_e k(1-k)I_0}{kT_e - T_v} \left[e^{-t/T_e} - e^{-kt/T_v} \right] \quad (7)$$

Where:

T_e is the simple beam time constant
 I_0 is the initial beam

Incorporating this expression into the beam decay model then gives a virtually perfect fit to beam decay. The standard deviation of the residual error is less than 0.15% which is the limit of accuracy of beam current measurements.

Storage Ring Operation

Having an accurate model, it is then possible to use it to gain some interesting insight into the operation of the Storage Ring.

Firstly it is desirable that the only significant process involved in beam decay should be gas scattering. Thus, in correct operation, other possible beam loss mechanisms such as betatron(-synchrotron) resonances, closed orbit errors or induced electromagnetic fields should only have a small effect. It should then be possible to establish a relationship between machine vacuum conditions and beam life, and any deviation from this relationship should indicate a machine malfunction. Unfortunately this has not proved to be so simple because of small long-term fluctuations in ion gauge calibration. Also the RF overvoltage has varied from time to time for operational reasons.

Nonetheless, a study of machine operating statistics over a period of six months has established a limit for the deviation of zero-beam lifetime from its expected value given by zero-beam pressure. Up to the time of writing all lifetimes which were outside this limit have been traced to a machine malfunction of some sort. So it is hoped that in the future the response of the storage ring diagnostics to abnormally short beam lifetimes will be considerably improved. A graph of zero-beam lifetime versus reciprocal zero-beam pressure together with the normal limits is shown in Figure 3.

Figure 3. Zero-beam life versus reciprocal zero-beam pressure

The low lifetimes are attributed to times when the horizontal and vertical betatron frequencies were allowed to approach one another too closely or when the beam was mis-aligned. The high lifetimes occurred when excessive RF power levels were used.

The average measured lifetime-pressure gradient is 11 hours-nanotorr¹⁾ compared with a theoretical value of 22 hours-nanotorr predicted by gas-scattering with the average RF overvoltage we are using. The most likely explanation of this is that the average pressure which is derived from the mean of many ion gauges located exclusively in straight sections is in error by that factor.

The most useful feature of the model is perhaps that it is now possible to identify the separate contributions to machine operating pressure of base pressure and beam induced gas desorption. This is particularly useful when machine behaviour is studied over a long period of time as is described in the next section.

Recovery from a Complete Vacuum Let-up

In November 1984 the whole machine was let up to atmospheric pressure while various modifications were carried out. A deliberate decision was taken not

to bake the vacuum system as was previous practice, but to run and outgas with synchrotron radiation. The history of zero-beam lifetime, zero-beam pressure and outgassing constant for the three months of machine operation at 1.8 GeV immediately following is shown in Figures 4, 5 and 6.

Figure 4. Zero beam life versus time

Figure 5. Zero beam pressure versus time

Figure 6. Outgassing constant versus time

From these graphs some immediate conclusions can be drawn. One is that the zero-beam pressure is exponentially reducing with an exponential time constant of about 130 days. However the average outgassing constant initially decayed much more rapidly with a time constant of about 32 days, but this rate slowed down considerably after the Christmas shutdown and presently the mean outgassing constant seems to be roughly steady at 25 picotorr/milliamper which gives an average lifetime of 7 hours for an injected beam of 300 mA. Any further decay appears to be very slow. Thus the high current lifetime is dominated by outgassing.

Individual values of the outgassing vary considerably from the mean. Much of this variation is attributed to the beam being steered to new positions to suit users. Other correlations have been established to shutdowns and operating energy changes.

General Conclusions

The beam decay model predictions have been compared with many actual beam decays and the correlation is at least as good as measurement error. Deviations from the model faithfully indicate machine malfunction or mis-adjustment. The operational history since the vacuum was let-up to atmospheric pressure indicates that, without a bake, the average high current beam lifetime eventually becomes dominated by beam induced outgassing which does not appear to reduce very quickly.

Reference

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CERN-TH.4461/86
ITP-UH 1/86

CONSTRAINTS ON SU(5)-TYPE LEPTOQUARKS

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A B S T R A C T

We analyze low energy effects of scalar leptoquarks in SU(5) unification schemes. Flavour changing couplings of leptoquarks with masses $M_{LQ} = 0$ (100 GeV) have to be strongly suppressed. Flavour diagonal couplings may be as large as $\lambda = O(1/10)$.

CERN-TH.4461/86
ITP-UH 1/86

June 1986

1. In the standard model of strong and electroweak interactions, and many of its variants, all scalars are colourless. More ambitious extensions (supersymmetry, substructure, grand unification, superstrings,...), however, imply a large variety of scalars with colour and other quantum numbers. Particularly interesting are 'leptoquarks' which couple to quark-lepton pairs and are colour triplets. These particles occur naturally in all unified models. With their unusual couplings they may cause new observable effects.

Leptoquarks (LQ) arising in SU(4) Pati-Salam unification [1], where lepton number is the fourth colour, have been analyzed previously [2]. Since they give rise to the highly suppressed decays $K_L^0 \rightarrow \mu e$, $K^+ \rightarrow \pi^+ \mu e$, either their mass must be at least 10 TeV (which is far beyond the range of energies accessible at present accelerators), or their couplings must be proportional to quark and lepton masses and hence highly suppressed.

On the other hand, an alternative kind of leptoquarks occurs in SU(5) unified theories. Since in SU(5) the (left-handed) fermions are in the $\underline{5}^*$ and $\underline{10}$ representations, scalars coupling to fermion pairs must be contained in the representations $\underline{5}$, $\underline{5}^*$, $\underline{10}$, $\underline{15}$, $\underline{45}$, $\underline{45}^*$ and $\underline{50}^*$. If the SU(5) symmetry is a remnant of an E_6 invariance where particles are grouped in 27 dimensional representations, a picture [3] suggested by superstring theories, leptoquarks occur only in $\underline{5}$ and $\underline{5}^*$ representations. The same is true for supersymmetric σ -models of quarks and leptons based on E_8 [4] whose particle spectrum is very similar to the one suggested by superstring theories.

The SU(2) x U(1) theory of electroweak interactions allows SU(2) singlet, doublet and triplet leptoquarks. Within the above pattern of symmetries, only the singlets are possible, and we will consider only those.

In the standard SU(5) model leptoquarks must be very heavy, $M_{LQ} \sim (10^{10} - 10^{15})$ GeV, in order to avoid too rapid proton decay. This is because the leptoquarks have also couplings to quark pairs, required by SU(5) invariance. Here we are interested in relatively light LQ's, with masses $M_{LQ} = O(M_W)$ which obviously requires that the couplings to quark pairs vanish, thus violating SU(5) invari-

ance^{*}). In this paper we investigate the phenomenological constraints on such light leptoquarks with non-SU(5) symmetric couplings.

An obvious lower bound on LQ masses of $M_{LQ} > 24$ GeV comes from e^+e^- -colliders because leptoquarks are charged. Stronger bounds follow from results at the CERN $p\bar{p}$ -collider where leptoquarks can be pair produced and decay with the typical signature of jets and leptons with large transverse momentum. Similar to scalar quarks a bound of 60 GeV can be given [6]. In this paper we concentrate on indirect effects of leptoquarks at low energies.

2. We start with the general couplings of isoscalar-colour triplet leptoquarks $G^{\alpha r}$ and antitriplets $G_{c\alpha}^r$ which are contained in the $\underline{5}$ and $\underline{5}^*$ representations of SU(5). Here α is the colour index and $r=1, \dots, n_G$ counts the number of LQ's. In the most popular models with E_6 invariance, as well as in the E_8 σ -models [4], $n_G=3$. The couplings read ($i, j=1, 2$ are SU(2) indices; $p, q, r=1 \dots n_G$ are 'flavour' indices):

$$\mathcal{L} = \bar{\lambda}_{pq}^r \varepsilon_{ij} q_p^{\alpha i} \ell_q^j G_{c\alpha}^r + \lambda_{pq}^r u_{c\alpha p} e_{cq} G^{\alpha r} + \text{c.c.}, \quad (1)$$

where $q^{\alpha i}$, ℓ_r^i are the left-handed quark and lepton doublets ($q_1^{\alpha 1} = u_L^\alpha, q_1^{\alpha 2} = d_L^\alpha, \dots$) and $u_{c\alpha r}$, e_{cr} are the complex conjugate right-handed up-quarks and charged leptons. In a specific model the couplings λ , $\bar{\lambda}$ are determined. Here we consider λ and $\bar{\lambda}$ as free parameters and study phenomenological constraints on these couplings.

The scalar fields $G^{\alpha s}$ and $(G_{c\alpha}^s)^* = G_c^{* \alpha s}$ have the same quantum numbers and can mix. In general, their mass matrix takes the form

$$\mathcal{M}^2 = \begin{pmatrix} M_{rs}^2 & M_{rs}^2 \\ \mu_{sr}^{*2} & \bar{M}_{rs}^2 \end{pmatrix}. \quad (2)$$

^{*}) This is a possibility in superstring theories [5]. We thank L. Ibañez for drawing our attention to this fact.

The effective four-fermion operator which is generated through the exchange of the mass eigenstates $\hat{G}^{QR} = U(G^{\alpha R}, G_c^{*\alpha R})$, $R=1\dots 2n_G$, is determined by the propagator matrix ξ satisfying

$$M^2 \xi = 1 M_{LQ}^2, \quad (3)$$

where M_{LQ} is an average leptoquark mass defined through

$$M_{LQ}^2 = (\det(M^2))^{\frac{1}{2n_G}}. \quad (4)$$

The matrix ξ is conveniently written as

$$\xi = \begin{pmatrix} \xi_{pq}^{(LL)} & \xi_{pq}^{(RL)} \\ \xi_{qp}^{(RL)*} & \xi_{pq}^{(RR)} \end{pmatrix}, \quad (5)$$

where the upper indices indicate the chiralities of the fermions connected by ξ .

For the four-fermion operator O , induced through leptoquark exchange, one obtains, after some Fierz-transformations

$$O = O_{LL} + O_{RR} + O_{RL}, \quad (6a)$$

$$O_{LL} = A_{pqrs}^{(LL)} (\bar{u}_{Lr} \gamma^\mu u_{Lp} \bar{e}_{Ls} \gamma_\mu e_{Lq} + \bar{d}_{Lr} \gamma^\mu d_{Lp} \bar{\nu}_{Ls} \gamma_\mu \nu_{Lq} - \bar{u}_{Lr} \gamma^\mu d_{Lp} \bar{e}_{Ls} \gamma_\mu \nu_{Lq} - \bar{d}_{Lr} \gamma^\mu u_{Lp} \bar{\nu}_{Ls} \gamma_\mu e_{Lq}), \quad (6b)$$

$$O_{RR} = A_{pqrs}^{(RR)} \bar{u}_{Rr} \gamma^\mu u_{Rp} \bar{e}_{Rs} \gamma_\mu e_{Rq}, \quad (6c)$$

$$O_{RL} = A_{pqrs}^{(RL)} (\bar{u}_{Rr} u_{Lp} \bar{e}_{Rs} e_{Lq} - \bar{u}_{Rr} d_{Lp} \bar{e}_{Rs} \nu_{Lq} + \bar{u}_{Rr} e_{Lq} \bar{e}_{Rs} u_{Lp} - \bar{e}_{Rs} d_{Lp} \bar{u}_{Rr} \nu_{Lq} + c.c.) , \quad (6d)$$

where

$$A_{pqrs}^{(LL)} = \frac{1}{2M_{LQ}^2} \bar{\lambda}_{pq} \xi^{(LL)} \bar{\lambda}_{rs}^* , \quad A_{pqrs}^{(RR)} = \frac{1}{2M^2} \lambda_{pq}^* \xi^{(RR)} \lambda_{rs} ,$$

$$A_{pqrs}^{(RL)} = \frac{1}{2M_{LQ}^2} \bar{\lambda}_{pq} \xi^{(RL)} \lambda_{rs} , \quad A_{pqrs}^{(LR)} = A_{rspq}^{(RL)*} . \quad (6e)$$

We will frequently use the notation

$$\overset{(-)}{\lambda}_{11} = \overset{(-)}{\lambda}_{ue} , \quad \overset{(-)}{\lambda}_{12} = \overset{(-)}{\lambda}_{u\mu} , \quad \overset{(-)}{\lambda}_{22} = \overset{(-)}{\lambda}_{c\mu} , \quad \dots \quad (7)$$

For up-type quarks we assume that mass and weak eigenstates coincide which can always be achieved by redefining the couplings $\lambda, \bar{\lambda}$.

The operators (6) contribute to various semileptonic processes. We note that the usual charged and neutral weak interactions also assume the form (6b), (6c). For the operator of ordinary β -decay one has, for instance, $A_{ueue}^{(LL)} = 2\sqrt{2}G_F \cos \theta_c$.

Let us consider first the decay of a pseudoscalar meson $M=(\bar{q}_1 q_2)$ into two leptons ℓ_1 and $\bar{\ell}_2$. The relevant hadronic matrix elements read

$$\langle 0 | \bar{q}_1(x) \gamma_\mu \gamma_5 q_2(x) | M(p) \rangle = i f_M p_\mu \varphi_M(x) , \quad (8a)$$

$$\langle 0 | \bar{q}_1(x) i \gamma_5 q_2(x) | M(p) \rangle = \frac{f_M m_M^2}{m_{q_1} + m_{q_2}} \varphi_M(x) , \quad (8b)$$

where p_μ , f_M , m_M and $\phi_M(x)$ are momentum, decay constant, mass and wave function of the meson M , and m_{q_i} are the current masses of the quarks q_i . From eqs. (6) and (8) one obtains for the decay width of $M \rightarrow \ell_1 \bar{\ell}_2$:

$$\Gamma = \frac{f_M^2 m_M^3}{64\pi} \lambda^{1/2} \left(1, \frac{m_{\ell_1}^2}{m_M^2}, \frac{m_{\ell_2}^2}{m_M^2}\right) \times \left\{ (|A_1|^2 + |A_2|^2) \left(1 - \frac{m_{\ell_1}^2 + m_{\ell_2}^2}{m_M^2}\right) + 4 \operatorname{Re}(A_1 A_2^*) \frac{m_{\ell_1} m_{\ell_2}}{m_M^2} \right\}, \quad (9a)$$

where

$$\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc, \quad (9b)$$

$$A_1 = \frac{m_M}{m_{q_1} + m_{q_2}} A^{(RL)} - \frac{m_{\ell_1}}{m_M} A^{(LL)} - \frac{m_{\ell_2}}{m_M} A^{(RR)}, \quad (9c)$$

$$A_2 = \frac{m_M}{m_{q_1} + m_{q_2}} A^{(LR)} - \frac{m_{\ell_2}}{m_M} A^{(LL)} - \frac{m_{\ell_1}}{m_M} A^{(RR)}. \quad (9d)$$

Contrary to leptonic meson decays mediated by weak interactions, eq. (9) contains terms proportional to $|A^{(RL)}|^2$, $|A^{(LR)}|^2$ which are not helicity suppressed, but instead enhanced by the mass ratio $\left(\frac{m_M}{m_{q_1} + m_{q_2}}\right)^2$. Note, however, that the coefficients $A^{(RL)}$, $A^{(LR)}$ are non zero only if the leptoquarks G^α and $G_C^{*\alpha}$ mix, i.e., if $\mu_{rs}^2 \neq 0$.

Important constraints on the couplings $\lambda, \bar{\lambda}$ can also be derived from μ -e conversion in matter. Following the work of Cahn and Harari [7] we compare, in the limit of zero incident momentum, the rate of ordinary muon capture,

$$\sigma \approx \frac{1}{16\pi} m_\mu^2 (g_V^2 + 3g_A^2) \frac{1}{\Lambda_F^4} , \quad (10)$$

derived from the operator

$$\begin{aligned} O_{CC} &= \frac{1}{2\Lambda_F^2} \bar{\nu}_L \gamma^\mu \mu_L \bar{n} \gamma_\mu (g_V - g_A \gamma_5) p , \\ \Lambda_F^2 &= (2\sqrt{2} G_F)^{-1} = (174 \text{ GeV})^2 , \end{aligned} \quad (11)$$

with the rate for coherent μ -e conversion following from the operator (6):

$$(\sigma)_{\mu e} = \frac{9}{32\pi} Z^2 m_\mu^2 \left(|A^{(LL)} + A^{(LR)}|^2 + |A^{(RR)} + A^{(RL)}|^2 \right) . \quad (12)$$

Here we have assumed a nucleus with equal number of protons and neutrons, and have used free Dirac spinors for electron and muon. Taking the nuclear form factor and the Pauli-exclusion principle in the case of μ -capture into account [8] we obtain for the ratio of cross sections

$$\frac{\sigma(\mu N \rightarrow e N)}{\sigma(\mu N \rightarrow \nu N)} = \frac{216 \Lambda_F^4}{(g_V^2 + 3g_A^2)} \left(|A^{(LL)} + A^{(LR)}|^2 + |A^{(RR)} + A^{(RL)}|^2 \right) . \quad (13)$$

3. We turn now to numerical results. Effects of leptoquark exchanges can appear in two ways. If the operator O of eq. (6) contributes to a process which already occurs through ordinary electroweak interactions, the new particles cause deviations from standard model predictions which are enhanced through interference with electroweak amplitudes. Otherwise the exchange of leptoquarks gives rise to new effects.

Consider first the former case, specifically corrections to β -decay [9]. In the interference, the coefficient $A_{ueue}^{(LL)} = A^{(LL)}$ dominates, yielding the modified decay rate

$$\Gamma \approx \Gamma_0 \left(1 + \frac{2 \operatorname{Re} A^{(LL)}}{2\sqrt{2} G_F} \right) \quad (14)$$

This leads to a deviation from Cabibbo-Kobayashi-Maskawa universality, since there is no contribution to the purely leptonic μ -decay. Assuming a typical allowed deviation from universality of 10^{-2} one obtains $\operatorname{Re} A^{(LL)} \leq 0.01/2\sqrt{2} G_F^2$ which, together with (6e) yields

$$M_{LQ} > (1.7 \text{ TeV}) (\bar{\lambda}_{ue} \zeta^{(LL)} \bar{\lambda}_{ue}^*)^{1/2} \quad (15)$$

The most stringent bound from processes not occurring in the standard model comes from coherent μ -e conversion. The recent experimental limit $\Gamma(\mu N \rightarrow e N) / \Gamma(\mu N \rightarrow \nu N) < 1.6 \cdot 10^{-10}$ [10] yields, together with eq. (13) ($g_V=1$, $g_A=-1.25$):

$$M_{LQ} > (150 \text{ TeV}) \left(|\bar{\lambda}_{u\mu} \zeta^{(LL)} \bar{\lambda}_{ue}^* + \bar{\lambda}_{ue}^* \zeta^{(RL)} \lambda_{u\mu}^*|^2 + |\lambda_{u\mu}^* \zeta^{(RR)} \lambda_{ue} + \bar{\lambda}_{u\mu} \zeta^{(RL)} \lambda_{ue}|^2 \right)^{1/4} \left(\frac{1.6 \cdot 10^{-10}}{\operatorname{BR}(\mu N \rightarrow e N)} \right)^{1/4} \quad (16)$$

We turn next to pseudoscalar meson decays for which precise experimental data exist. Clearly the contribution of the chirality changing operator O_{RI} dominates the helicity suppressed terms of O_{LL} and O_{RR} (unless $A^{(RL)} \ll 1$). Neglecting corrections of order m_l^2/m_M^2 one obtains

$$\Gamma(M \rightarrow l, \bar{l}) = \frac{f_M^2 m_M^3}{32 \pi} \left(\frac{m_M}{m_{q_1} + m_{q_2}} \right)^2 |A^{(RL)}|^2 \quad (17)$$

Applying (17) to the decays $\pi^0 \rightarrow \mu^\pm e^\mp$, $\pi^0 \rightarrow e^+ e^-$, $\pi^\pm \rightarrow e^\pm \nu$ yields the bounds (cf. eq. (6e)):

$$M_{LQ} > (410 \text{ GeV}) \left(|\bar{\lambda}_{u\mu} \xi^{(RL)} \lambda_{ue}|^2 + |\bar{\lambda}_{ue} \xi^{(RL)} \lambda_{u\mu}|^2 \right)^{1/4} \cdot \left(\frac{7 \cdot 10^{-8}}{\text{BR}(\pi^0 \rightarrow \mu^\pm e^\mp)} \right)^{1/4}, \quad (18a)$$

$$M_{LQ} > (320 \text{ GeV}) |\bar{\lambda}_{ue} \xi^{(RL)} \lambda_{ue}|^{1/2} \left(\frac{1.8 \cdot 10^{-7}}{\text{BR}(\pi^0 \rightarrow e^+ e^-)} \right)^{1/4}, \quad (18b)$$

$$M_{LQ} > (8.8 \text{ TeV}) \left(\sum_i |\bar{\lambda}_{ue_i} \xi^{(RL)} \lambda_{ue}|^2 \right)^{1/4} \left(\frac{1.2 \cdot 10^{-4}}{\text{BR}(\pi^\pm \rightarrow e^\pm \nu)} \right)^{1/4}. \quad (18c)$$

Here we have normalized the branching ratios to the experimental bounds, or the measured values, which are in agreement with the standard electroweak theory.

For comparison, in the case $A^{(RL)} = 0$ one obtains for the helicity suppressed amplitudes $A^{(LL)}$

$$M_{LQ} > (90 \text{ GeV}) |\bar{\lambda}_{u\mu} \xi^{(LL)} \bar{\lambda}_{ue}^*|^{1/2} \left(\frac{7 \cdot 10^{-8}}{\text{BR}(\pi^0 \rightarrow \mu^\pm e^\mp)} \right)^{1/4}, \quad (19a)$$

$$M_{LQ} > (5 \text{ GeV}) (\bar{\lambda}_{ue} \xi^{(LL)} \lambda_{ue}^*)^{1/2} \left(\frac{1.8 \cdot 10^{-7}}{\text{BR}(\pi^0 \rightarrow e^+ e^-)} \right)^{1/4}, \quad (19b)$$

$$M_{LQ} > (120 \text{ GeV}) \left(\sum_i |\bar{\lambda}_{ue_i} \xi^{(LL)} \bar{\lambda}_{ue}^*|^2 \right)^{1/4} \left(\frac{1.2 \cdot 10^{-4}}{\text{BR}(\pi^\pm \rightarrow e^\pm \nu)} \right)^{1/4}. \quad (19c)$$

A stringent bound follows from the decay $\eta \rightarrow \mu^+ \mu^-$, which is particularly enhanced by the mass ratio $\left(\frac{m_M}{m_{q_1} + m_{q_2}} \right)^2$ [11]

$$M_{LQ} > (180 \text{ GeV}) |\bar{\lambda}_{u\mu} \xi^{(RL)} \lambda_{u\mu}|^{1/2} \left(\frac{6.5 \cdot 10^{-6}}{\text{BR}(\eta \rightarrow \mu^+ \mu^-)} \right)^{1/4}. \quad (20)$$

Coming next to the K-decays we first notice that SU(5)-type leptoquarks do not give rise to decays into charged leptons, such as $K_L^0 \rightarrow \mu^+ \mu^-$, $K_L^0 \rightarrow \mu^\pm e^\mp$ etc. Hence the strong constraints which apply to Pati-Salam type leptoquarks are avoided. Severe bounds can be derived, however, from leptonic and semi-leptonic decays involving neutrinos. From $K^+ \rightarrow e^+ \nu$ we find [11]

$$M_{LQ} > (15 \text{ TeV}) \left(\sum_i |\bar{\lambda}_{ue} \xi^{(u)} \bar{\lambda}_{ce}^*|^2 \right)^{1/4} \left(\frac{1.5 \cdot 10^{-5}}{\text{BR}(K^+ \rightarrow e^+ \nu)} \right)^{1/4}, \quad (21)$$

whereas the limit on $K^+ \rightarrow \pi^+ \bar{\nu} \nu$ yields

$$M_{LQ} > (7.3 \text{ TeV}) \left(\sum_{ij} |\bar{\lambda}_{ue_i} \xi^{(u)} \bar{\lambda}_{ce_j}^*|^2 \right)^{1/4} \left(\frac{2.9 \cdot 10^{-6}}{\text{BR}(K^+ \rightarrow \pi^+ \bar{\nu} \nu)} \right)^{1/4}. \quad (22)$$

For simplicity, we have considered in (22) only the contribution of O_{LL} since there is no helicity suppression in this case.

Finally we turn to D-decays where recently new experimental results have been obtained for the leptonic decay modes for $D^0 \rightarrow \mu^+ \mu^-$ [12] and $D^0 \rightarrow \mu^\pm e^\mp$ [13]. The corresponding bounds read

$$M_{LQ} > (1.8 \text{ TeV}) |\bar{\lambda}_{u\mu} \xi^{(u)} \lambda_{u\mu}|^{1/2} \left(\frac{1.0 \cdot 10^{-5}}{\text{BR}(D^0 \rightarrow \mu^+ \mu^-)} \right)^{1/4}, \quad (23a)$$

$$M_{LQ} > (520 \text{ GeV}) \left(|\bar{\lambda}_{u\mu} \xi^{(u)} \lambda_{ce}|^2 + |\bar{\lambda}_{ue} \xi^{(u)} \lambda_{c\mu}|^2 \right)^{1/4} \cdot \left(\frac{1.5 \cdot 10^{-3}}{\text{BR}(D^0 \rightarrow \mu^\pm e^\mp)} \right)^{1/4}. \quad (23b)$$

We note that due to the SU(2) invariance of the effective four-fermion operator O in eq.(6) K- and D-decays are related. The K-decays $K^+ \rightarrow \bar{e}_i \nu_j$ or $K^+ \rightarrow \pi^+ \bar{\nu}_i \nu_j$, for instance, would imply the D-decays $D^0 \rightarrow \bar{e}_i e_j$.

4. The bounds we have derived depend crucially on the coefficients $\lambda, \bar{\lambda}$ which measure the coupling strength of leptoquarks to quark-lepton pairs. If the λ 's are of the order of usual gauge couplings, i.e., $\lambda \approx \frac{1}{2}$, the mass values given in section 3 are reduced by a factor of 2. In this case, unless the couplings $\lambda, \bar{\lambda}$ have a special flavour structure, the leptoquark masses have to be of the order of 100 TeV. On the contrary, mass suppressed couplings of the form

$$\lambda_{q\ell} \approx \frac{1}{2} \frac{(m_q m_\ell)^{1/2}}{m_W} \quad (24)$$

would allow leptoquark masses of order 100 GeV.

As the coefficients $\lambda, \bar{\lambda}$ are Yukawa couplings, a particular flavour structure is indeed expected. Flavour diagonal couplings of mass eigenstates, for instance, which appear natural since leptoquarks and Higgs doublets are in the same SU(5) multiplets, eliminate the bounds (16), (18a), (20) and (23a). Leptoquarks with masses of order 100 GeV require, however, that also the couplings in (18b), (18c), (21) and (23) are suppressed. This would imply that the coefficients $\xi^{(RL)}$, and hence the mixing between $G^{\alpha r}$ and $G_c^{*\alpha r}$, are small. Finally, the bound (22) can only be avoided if quark-lepton pairs couple unsuppressed only to leptoquarks of the same generation! We conclude that unsuppressed leptoquark couplings are only possible if the theory has an extended 'Higgs sector' of scalar fields where the dominant couplings are flavour diagonal and where right- and left-handed fermions interact most strongly with their own leptoquarks. We emphasize that such a structure is indeed conceivable in E_6 -type N=1 supergravity theories or in the supersymmetric σ -models based on E_8 [4].

For the pattern of leptoquark couplings described above the constraint (14) from quark-lepton universality still remains. A leptoquark $G_{c\alpha}^1$ of the first generation with mass $M_{LQ} = O(100 \text{ GeV})$ requires a coupling $\bar{\lambda}_{ue}^1$ not larger than $O(\frac{1}{10})$. The coupling λ_{ue}^1 of $G^{\alpha 1}$ to right handed electron-up-quark pairs can even be stronger. These couplings are also compatible with constraints derived from the anomalous magnetic moment of the electron [14]. The contribution to $(g-2)_{e,\mu}$ due to virtual leptoquarks is suppressed

because of the small mixings between $G^{\alpha r}$ and $G_C^{*\alpha r}$.

As our analysis shows it is indeed possible that within an extended 'Higgs sector' of scalar fields, there exist colour-triplet scalars with masses of order 100 GeV which, together with SU(2) Higgs doublets, form complete $\underline{5}$ or $\underline{5}^*$ representations of SU(5). The couplings $\lambda, \bar{\lambda}$ of these colour triplets cannot be SU(5) symmetric in order to conserve baryon and lepton number. It is an interesting possibility that the diquark couplings vanish and that the colour-triplets consequently appear as leptoquarks. Flavour diagonal couplings of such leptoquarks do not have to be mass suppressed; they are, however, expected to be smaller than gauge couplings ($\bar{\lambda} < \frac{1}{10}$). These couplings are large enough to allow sizeable production cross sections at HERA.

We would like to thank R. Rückl, B. and F. Schrempp and P.M. Zerwas for discussions on leptoquark production at HERA. We have been informed by D.V. Nanopoulos that questions similar to those discussed in this paper have also been studied by B.A. Campbell et al. [15].

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