CHEEP

AN e-p FACILITY IN THE SPS

CHEEP Study Groups
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ABSTRACT

The addition of a 25 GeV (30 GeV) electron ring to the 400 GeV proton synchrotron (SPS) at CERN would substantially extend the kinematical region ($Q^2, \nu$) accessible to observation, compared to the present neutrino and muon experiments. This report presents the theoretical arguments in favour of such a project, together with a possible detection system and a design proposal. The studies indicate that e-p collisions at these energies would enable fundamental investigations to be made of the strong, electromagnetic, and weak interactions; in particular, the weak interactions would be probed, for the first time, in the region of their characteristic mass. Detailed calculations of the expected rates are presented, with a discussion of detection possibilities taking into account the anticipated background. A detector, conceived to cover a great variety of events whilst still fitting into the SPS tunnel, is proposed, and its functioning is explained in detail. Finally, a feasibility study of such a colliding-beam machine is presented. A 25 (30) GeV ring in the SPS tunnel would provide longitudinally polarized electrons or positrons colliding with the SPS protons during their acceleration for the fixed-target programme; the luminosity would be between $10^{31}$ and $10^{32}$ cm$^{-2}$ sec$^{-1}$, depending on the mode of operation of the SPS.
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*edited by J. Ellis*

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INTRODUCTION

B.H. Wiik
INTRODUCTION

It was proposed in late 1976 to install in the CERN SPS tunnel an electron storage ring capable of reaching 25 GeV (30 GeV), and to collide the electrons (positrons) with the protons in the main ring.

The kinematical region which can be exploited with this option is shown here. Also shown is the region which now is under investigation with the 300 GeV muon and neutrino beams available at the CERN SPS and at FNAL. This program can be extended by nearly two orders of magnitude in $Q^2$ and $\nu$ using CHEEP -- the CERN High-Energy Electron Proton Option. This extension would make it possible, for the first time, to explore the electromagnetic and the weak interaction in a region of $Q^2$ comparable to the characteristic mass of the weak interaction squared.

Early in 1977 an ECFA Study Group was formed, with a strong CERN participation. The aim of this group was to assess the physics interest of a high-energy electron-proton facility, identify possible experimental problems and investigate the detectors needed to exploit the physics, and to establish the technical feasibility of the scheme put forward. The present report is a summary of the work done in the Study Group.

The physics interest

The study of neutral and charged weak currents is one of the prime motivations for constructing a large electron-proton colliding beam facility. Present and foreseeable experiments only probe the weak interaction at values of $\sqrt{Q^2}$ which are small compared to the characteristic mass-scales of the weak interaction. Unitarity effects should become important at centre-of-mass energies of a few hundred GeV, while gauge theories suggest a unification with electromagnetism on a boson mass scale around 100 GeV. Measurements at CHEEP will extend weak interaction investigations into a new regime where $\sqrt{Q^2}$ is comparable with these scales. Measurements at such large values of $Q^2$ will presumably provide the answers to some crucial questions:

---

1) Measurements of $e + p \rightarrow e' + X$ with left- and right-handed electrons and positrons at large $Q^2$ will show unambiguously if the electromagnetic and the weak interaction are indeed manifestations of a single interaction (as believed at present) or not. The mechanism selected by nature to achieve the (conjectured) unification might be unravelled. In particular, these measurements afford precision probes of the current gauge theories.

2) Measurements of the charged weak current in $e + p \rightarrow \nu + X$ will establish if the weak interaction is damped by a single intermediate vector boson, or if a more complicated mechanism, perhaps involving several such bosons, is needed. Left-handed positrons and right-handed electrons are very sensitive probes for new weak currents.

It is of crucial importance to identify the mass spectrum of leptons and quarks. The large production cross-section and the spectacular signatures make CEEP an ideal machine for searching for new electron-like leptons with masses up to 100 GeV. New quarks can be produced singly in the weak interaction, and the rates are sufficient for exploring the quark spectrum up to a mass of about 60 GeV. Electroproduction is expected to be a fertile source of associated production of new flavours. Detailed computations indicate that this mechanism will permit the quark spectrum to be explored for quark masses up to 20-30 GeV.

The well-known properties of the electromagnetic interaction can be used in the reaction $e + p \rightarrow e' + X$ to probe the structure of the strong interactions. Such measurements at SLAC gave the first hints that the charge in the proton is not uniformly distributed but is tied to point-like objects, the partons.

With CEEP, the present experiments can be extended by nearly two orders of magnitude in $Q^2$ down to distances of $10^{-16}$ cm. Such measurements are obviously of fundamental importance for our understanding of the strong interaction. In particular, a precise study of scaling violations might reveal the structure of the strong interactions — especially whether present ideas about a gauge theory of the strong interactions are justified or not. Stringent tests of this picture are obtained from measurements of the final state; for example, if our present understanding is correct, then not only "quark" but also "gluon" jets will be observed.

Photoproduction at very high energies can be studied with CEEP in great detail, since millions of events are produced per day with $Q^2 < m_n^2$. If the photons continue to behave like hadrons at high energies, this allows the study of all conventional hadron physics with the advantage that the beam is polarized and the variation with $Q^2$ can be measured. In addition, a photon beam is a copious source of the diffractive production of heavy vector mesons, and the Primakoff effect gives a large production cross-section for all particles with a sizeable coupling to two photons.

Experiments

The first task of the experimental working group was to establish the feasibility of performing deep inelastic scattering experiments using the proposed facility.

The circulating proton beam is a potentially serious source of background via beam-gas scattering or other loss mechanisms. To evaluate this problem, measurements were made of the background in the SPS tunnel for both stored and accelerated beams. The rates were nearly the same for both running conditions and were found to be easily manageable when scaled to CEEP running conditions. The same conclusion was also reached for the background
associated with the electron beam. The synchrotron radiation from the high-energy electron beam represents a more serious problem. However, detailed considerations indicate that this background is not prohibitive and can be overcome by a careful layout of the interaction region and the adjacent machine elements.

Based on present data, deep inelastic charged and neutral current events will in general lead to final states with a three-jet structure: a lepton, new or old, travelling at a large angle with respect to the beam line with its transverse momentum balanced by the particles in the current jet; the proton jet, resulting from the break-up of the proton, will travel along the initial proton direction with small net transverse momentum.

A calorimeter-type device which covers a large solid angle and is capable of detecting both electrons and hadrons was investigated and found to be suitable for a large class of deep inelastic inclusive experiments. Such a detector is also suitable for charged current events, provided that hadrons which travel at small angles with respect to the proton beam are measured.

Particles in the lepton jet or in the current jet will in general emerge at large angles and can therefore be identified and measured. This makes it feasible to search for new leptons and new quarks with a high efficiency. The detectors that are suitable for such experiments are similar to those discussed for large e^+e^- machines, and are in general based on solenoids and toroids.

Photoproduction requires special detectors with good electron detection at small angles with respect to the electron beam.

The machine

The proposal calls for installing an electron storage ring above the synchrotron in the SPS tunnel and for colliding the electrons in this ring with the protons in the main ring. Collisions can take place during acceleration for a proton energy range between 145 GeV and 400 GeV. This makes it possible to carry out at least a part of the normal synchrotron program in parallel with the e-p program. Alternatively, collisions can take place at a fixed proton energy during an intermediate plateau in the SPS cycle. The short proton storage time in a fast-cycling accelerator such as the SPS alleviates all problems associated with the lifetime of the proton beam.

The peak luminosity per interaction region is $0.5 \times 10^{32}$ cm$^{-2}$ sec$^{-1}$ ($0.75 \times 10^{32}$ cm$^{-2}$ sec$^{-1}$ without slow proton extraction) for 270 GeV protons on 25 GeV electrons. This luminosity, even when reduced by an order of magnitude, is still sufficient to explore a region in the $Q^2$,x diagram well above the one attainable at the SPS or at FNAL. If CHEEP could operate as a dedicated facility, luminosities in excess of $10^{32}$ cm$^{-2}$ sec$^{-1}$ would be achieved.

The working group has investigated in great detail the possibility of obtaining longitudinally polarized beams of electrons and positrons. All work to date shows that the electrons (positrons) will indeed be transversely polarized in CHEEP with a build-up time of 0.5 h over a wide range of energies. In order to fully exploit this feature, the layout of the magnetic elements in the interaction region provides rotation of the transverse polarization into the longitudinal direction over a wide range of electron energies.
The machine study group concentrated on a number of points which appeared to be of crucial importance for establishing the basic feasibility of the proposed scheme. In order to obtain a realistic assessment of the impact on the fixed-target program, the implications for the operation of the SPS were studied. Great attention was paid to the layout of the insertion which determines performance and flexibility to a great extent. All the principal parameters for the insertion and for the electron ring were determined, providing a solid basis for a full design report.

Concluding remarks

The present report shows that crucial information on strong, electromagnetic, and weak interactions can be obtained from experiments with longitudinally polarized electrons and positrons colliding with unpolarized protons at high energies. Although new machines must always be justified by extrapolating present knowledge, it seems very likely that entirely new phenomena will be found in the large kinematical area opened up by CHEEP. We can only speculate whether unexpected new interactions or particles will be found. However, the large momentum transfers available with CHEEP give the possibility of exotica such as free quarks and gluons, monopoles, Higgs bosons, and intermediate vector bosons, if such things exist in energy ranges an order of magnitude larger than those explored so far.

The study further demonstrates the feasibility of performing the experiments and it lists the requirements which the detectors must fulfil. The experimental program is rich, and it can only be exploited by using several different types of detectors. The report, although incomplete, demonstrates the feasibility of the proposed scheme. The luminosity and the c.m.s. energies are sufficiently large to carry out the physics program discussed. However, more work is needed to bring the study up to the level required for a full design report.
CHAPTER I
THEORETICAL REMARKS ABOUT e-p COLLISIONS

Summary, prepared by J. Ellis,
of work by a theoretical group including:

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1. **INTRODUCTION**

As emphasized in the general introduction\(^1\), the kinematic range of e-\(p\) collisions opened up by CEEP should enable fundamental investigations of the strong, weak, and electromagnetic interactions to be made in a new régime of \(Q^2\) and \(\nu\). On the one hand, the long lever arm provided by the large momentum transfers offered by CEEP should enable us to see whether the strong interactions can be described by a field theory\(^2\)\(^-\)\(^4\). If so, we should be able to distinguish different models, and in particular refute or confirm the predictions of asymptotically free gauge theories (QCD)\(^5\). As far as the weak interactions are concerned, the momentum transfer range offered by CEEP is comparable with the mass-scale \(\mathcal{O}(100)\) GeV characteristic of the weak interactions, where we may expect deviations from the four-fermion theory to occur. In particular, many theories\(^6\) anticipate intermediate vector bosons \(W^\pm, Z^0\), which have masses in the range (50-100) GeV. Neutral and charged weak interactions are comparable to electromagnetism, interference effects are \(\mathcal{O}(1)\), and any damping or \(W^\pm\) and \(Z^0\) propagator effects should be clearly visible. Competing unified theories of the weak and electromagnetic interactions can therefore be distinguished. Deep inelastic collisions should also be a copious source of the new quarks and leptons expected in many models.

Section 2 of this chapter examines some aspects of strong interaction studies accessible with CEEP:

- Tests of scaling\(^7\), deviations from it, and the qualitative predictions of different field theories of the strong interactions\(^1\)-\(^5\).
- The pair-production of new quarks and their kinematic signatures.
- Studies of the hadronic final state, searches for jets, \(p_T\) broadening and the like.
- Photoproduction studies at low \(Q^2\), such as the production of new mesons.

Section 3 examines some aspects of weak interaction studies accessible with CEEP:

- Neutral current effects and the possibilities for distinguishing different models\(^4\).
- Charged current effects, including the sensitivity to deviations from the four-fermion interaction and possible \(W^\pm\) propagator effects.
- The search for new weakly interacting particles, such as heavy leptons, new quarks, and intermediate vector bosons.

The list of topics studied is not intended to be exhaustive, but rather a sampling of what future prospects intrigue us most at the present time. We may hope that other, more exciting, surprises lie in store. Since the large momentum transfers provided by CEEP guarantee access to short distances \(\mathcal{O}(10^{-16} \text{ cm})\), we feel that this machine is particularly susceptible to the unanticipated.

2. **STRONG INTERACTION STUDIES**

The approximate scaling discovered in deep inelastic scattering\(^4\) of electrons, neutrinos, and muons made us look at strong interactions in a new way. It now seems that hadrons contain quark-parton constituents which are rather point-like and carry the weak and electromagnetic couplings of strongly interacting matter\(^3\). The nucleon seems to contain other
constituents without these couplings, which one may call "gluons" without any prejudice as to their nature. These quark and "gluon" constituents are now used in models of many aspects of strong interactions, including spectroscopy, large p_T phenomena, the production of lepton pairs, and multiparticle production. They are also essential ingredients in our understanding of the possible unification of weak and electromagnetic interactions. It is clearly very important to understand more about these constituents, studying in particular:

- the continued validity down to distances \(O(10^{-16} \text{ cm})\) of the concept of an essentially point-like constituent;
- the interactions between these constituents and their possible description by a field theory such as an asymptotically free gauge theory\(^7\); \(^5\);
- the way that partons are bound into hadrons, and the way they recombine into jets of hadrons after a hard collision\(^9\); and
- the existence of possible new types of constituent.

We will see that CHEEP is the ideal way to study the first two of these questions, is of comparable value to a high-energy e+e- machine for studying recombination and jets, and has useful capabilities in the study of new types of quark.

2.1 Scaling and the alternatives

In the one-photon exchange approximation the cross-section for e + p \(\rightarrow e' + X\) can be written in terms of two dimensionless structure functions:

\[
\frac{d^2\sigma^{\gamma^*Y}}{dx dy} = \frac{4\pi a^2}{s x^2 y^2} \left\{ (1-y) F_2(x,Q^2) + y^2 x F_1(x,Q^2) \right\} ,
\]

where the kinematics and definitions used are discussed in detail in Chapter II. If \(F_1\) and \(F_2\) scale in the manner suggested by Bjorken\(^7\) and interpreted via the parton model\(^9\) of Fig. I.1,

\[
F_{1,2}(x,Q^2) = F_{1,2}(x) ,
\]

then there will be usable event rates for studying e + p \(\rightarrow e' + X\) in most of the kinematic range opened up by CHEEP. If the partons coupled to the photon have spin 1/2 (quarks, for example), then \(F_1(x)\) and \(F_2(x)\) are related by

\[
F_2(x) = 2x F_1(x) ,
\]

as was shown by Callan and Gross\(^10\). Equation (I.3) means that in the scaling limit

\[
R \equiv \sigma_L/\sigma_T \rightarrow 0 .
\]

Figure I.2 shows the rates calculated\(^2\) using a naive scaling quark-parton model\(^11\) in the one-photon approximation (I.1), showing that there will be ample opportunity to study the dynamics of quark-partons and look at any deviations from the naive point-like picture. Figure I.3 shows event rates for the charged current reaction e + p \(\rightarrow \nu + X\), under various assumptions. Again we find usable event rates at large \(Q^2 = O(10^7)\).
Exact scaling in the form (1.2) proposed by Bjorken\(^7\) does not seem to hold in present electron, neutrino, and muon scattering data. This may just be a threshold phenomenon, but it is widely believed that the scaling violations may extend out to large \(Q^2\). Exact scaling seems to be unobtainable in the context of quantum field theory\(^{11}\). Some theoreticians therefore abandon the field theoretical framework and retain scaling\(^{14}\), while others study the scaling violations found in different field theories\(^5\). It may well be that both approaches are wrong, and quarks are themselves composed of more fundamental constituents\(^{15}\) (preons? prequarks? straton?) which will show up at distances shorter than those accessible to present experiments. We now examine these alternatives in some detail.

2.1.1 Field theories of the strong interactions?

Every known field theory violates scaling\(^{11}\). There is one model known -- an asymptotically free gauge theory (QCD)\(^5\) -- where the violations are logarithmic in \(Q^2\). Other theories are believed\(^1\) to have power-law \(Q^2\) variations. However, the qualitative features of the underlying physical mechanisms and patterns of scaling violation are somewhat model-independent.

Consider\(^5\) a low \(Q^2\) virtual photon striking a quark-parton constituent carrying a fraction \(x\) of the longitudinal momentum of the proton as indicated in Fig. I.4. As the \(Q^2\) of the virtual photon is increased, its spatial resolving power is improved, and sometimes it may separate the large "low \(Q^2\)" quark into a smaller quark and accompanying "gluon", as shown in Fig. I.4b. This can just be regarded as gluon bremsstrahlung, or a semistrong analogue of Compton scattering. The momentum fraction \(x'\) of the quark seen at high \(Q^2\) is clearly lower on the average than the \(x\) of the low \(Q^2\) quark. Another mechanism for changing the structure function as \(Q^2\) increases is indicated in Fig. I.5. A low resolving power, low \(Q^2\) photon fails to interact with a "gluon" because it has no electromagnetic charge. A high \(Q^2\) virtual photon may resolve the "gluon" into a quark-antiquark pair of constituents, as illustrated in Fig. I.5b. This process can be regarded as quark pair creation in the strong "gluon" field of the nucleon. Again, the momentum fraction \(x'\) carried by the quark or antiquark struck at high \(Q^2\) will in general be lower than the \(x\) of the low \(Q^2\) gluon. Therefore the effect of these lowest-order strong "radiative corrections" at large \(Q^2\) is to deplete the structure functions \(F_1(x,Q^2)\) at large \(x\), and to build them up at low \(x\), as indicated in Fig. I.6.

This intuitive lowest-order picture can be made more mathematically rigorous by considering the moments of the structure functions \(F_1(x,Q^2)\):

\[
\int_0^1 dx \ x^{n-1} F_1(x,Q^2) = \nu_{1}^{n}(Q^2) \ . \quad (I.5)
\]

Exact scaling would require that the \(\nu_{1}(Q^2)\) be constant as \(Q^2 \to \infty\). It has been shown\(^1\) from the positivity of the structure functions that

\[
\frac{\nu_{1}^{n+m}(Q^2)}{\nu_{1}^{n}(Q^2)} \leq \text{some constant } C^{\text{RAM}}, \quad \text{for any } n \geq 2, \quad m > 0 \ . \quad (I.6)
\]
while $\mu_1^2(Q^2) \to \text{constant}$. If the structure functions do not behave pathologically, Eqs. (1.5) and (1.6) mean that they go to zero faster as $x \to 1$, while the area under $F_2(x,Q^2)$ is asymptotically constant as $Q^2 \to \infty$, in qualitative agreement with the intuitive picture\textsuperscript{16) of Fig. 1.6.

These qualitative arguments were independent of the details of the behaviour of the strong interaction coupling constant strength (or strengths) as $Q^2 \to \infty$ (short distances). Various alternatives can be imagined\textsuperscript{15):} perhaps $g \to 0$ (asymptotic freedom), or $g \to \infty$ for some fixed value $g^* \neq 0$, or something else happens. This last case cannot yet be analysed in general, and we can only refer back to the previous qualitative arguments. In the case $g \to g^* \neq 0$, the moments (1.5) violate scaling by powers of $Q^2$ \textsuperscript{16):

$$
\mu_1^3(Q^2) \sim \left( \frac{Q^2}{\Lambda^2} \right) a_1^3, \quad a_1^m = 0 \text{, } a_1^{n+m} \sim a_1^n \text{ for } m > 0 .
$$

(1.7)

A realistic fixed-point theory of this type with a calculable value of $g^*^2$ is not known. However, the forms of the $a_1^n$ for two theories which do not have $g^* = 0$ are\textsuperscript{16)}:

$$
\begin{aligned}
\text{scalar} & \quad a_1^n = \frac{g_s^2}{24\pi^2} \left( 1 - \frac{2}{n(n+1)} \right) \\
\text{gluons} & \quad a_1^n = \frac{g_s^2}{12\pi^2} \left( 1 - \frac{2}{n(n+1)} + 4 \sum_{j=2}^n \frac{1}{j} \right).
\end{aligned}
$$

(1.8)

Abelian vector

$$
\begin{aligned}
\text{gluons} & \quad a_1^n = \frac{g_v^2}{16\pi^2} \left[ 1 - \frac{2}{n(n+1)} + 4 \sum_{j=1}^n \frac{1}{j} \right].
\end{aligned}
$$

(1.9)

The forms (1.8) and (1.9) are appropriate to flavour non-singlet (valence quark) combinations of structure functions, which are expected to control the behaviour of $vW^2_{\text{up}}$ for $x \geq 0.4$. Predictions\textsuperscript{19}) based on (1.8) and (1.9) are shown in Fig. 1.7: the fixed-point $g_v^*$ was chosen so as to mimic as far as possible the predictions of asymptotically free gauge theory in the range $2 \leq Q^2 \leq 30$.

The only known theory in which $g \to 0$ as $Q^2 \to \infty$ is a non-Abelian colour gauge theory (QCD)\textsuperscript{19).} In this case the $Q^2$ dependence of the moments is logarithmic and exactly known for the combinations of structure functions which can be interpreted in terms of valence quark distributions:

$$
\mu_1^3(Q^2) \sim \left( \ln \frac{Q^2}{\Lambda^2} \right)^3 \delta^n \left[ 1 + \mathcal{O} \left( \ln \frac{\ln Q^2/\Lambda^2}{\ln Q^2/\Lambda^2} \right) \mu_1^3 \right].
$$

(1.10)

where

$$
\delta^n = \frac{4}{35-2F} \left[ 1 - \frac{2}{n(n+1)} + 4 \sum_{j=2}^n \frac{1}{j} \right].
$$

(1.11)
For non-valence combinations of structure functions

\[
\mu_n^\perp(Q^2) \approx \left( \ln \frac{Q^2}{\Lambda^2} \right)^{-6n} \left[ 1 + \mathcal{O} \left( \ln \frac{\ln Q^2}{\ln Q^2/\Lambda^2} \right) \right] \nu_n^\perp
\]

(1.12)

where the $d_{\perp}^n$ are eigenvectors of the matrix

\[
\begin{pmatrix}
\frac{4}{33-2F} \left( 1 - \frac{2}{n(n+1)} + 4 \sum_{j=2}^{n} \frac{1}{j} \right) & -\frac{4}{35-2F} \frac{2(n^2+n+2)}{n(n-1)} \\
-\frac{F}{22-\frac{15F}{4}} \frac{4(n^2+n+2)}{n(n+1)(n+2)} & \frac{9}{33-2F} \left( \frac{1}{3} - \frac{4}{n-1} - \frac{4}{n(n+1)(n+2)} + 4 \sum_{j=2}^{n} \frac{1}{j} \right)
\end{pmatrix}
\]

(1.13)

In the above formulae, $F$ is the number of quark flavours and $\Lambda^2$ is an \textit{a priori} unknown parameter which must be determined phenomenologically. The $\mathcal{O}(\ln \ln Q^2/\Lambda^2/(\ln Q^2/\Lambda^2))$ correction terms in (1.10) have been calculated\(^{26}\), and are known to have little effect on phenomenological analyses made using just the leading-order terms\(^{21}\). QCD is widely regarded as the leading candidate for a field theory of the strong interactions: the formulae (1.9) to (1.13) are unambiguous predictions by which it stands or falls.

Under suitable technical assumptions, the moments (1.5) can be inverted, and the structure functions $F_i(x,Q^2)$ expressed in terms of the structure functions $F_i(x,Q^2)$ measured at some base-line $Q^2$. This evolution of the structure functions can be expressed in terms of $Q^2$-dependent effective quark distributions $q(x,Q^2)$ and this has been done by several authors\(^{16,22}\). For practical purposes it is sometimes useful to have a compact approximation to the $Q^2$ dependence of the quark distributions, and one such parametrisation for valence quarks is\(^{22}\) [the quark distributions, $q_i = d_V, u_V$, etc., are defined so that $F_2(x) = 2xP_1(x) = x^2Q^2\frac{d}{dx}(q_1(x))$]:

\[
\begin{align*}
xd_V(x,Q^2) &= \frac{1}{B(a_1, a_2+1)} x^{a_1}(1-x)^{a_2} \\
 x(d_V(x,Q^2) + u_V(x,Q^2)) &= \frac{3}{B(c_1, a_4+1)} x^{a_3}(1-x)^{a_4}
\end{align*}
\]

(1.14)

where $B(a,b)$ is the usual beta function, and

\[
\begin{align*}
\alpha_1 &= 0.97 - 0.27 s & \alpha_2 &= 3.55 + 0.81 s \\
\alpha_3 &= 0.72 - 0.19 s & \alpha_4 &= 2.8 + 0.81 s \\
s \equiv \ln \left( \frac{\ln Q^2/\Lambda^2}{\ln Q_0^2/\Lambda^2} \right), \quad Q_0^2 = 2 \text{ GeV}^2 .
\end{align*}
\]

(1.15)
Buras and Gaemers also give parametrizations of sea quark and antiquark distributions. These formulae describe present electron and muon production data rather well, if one takes \( \Lambda \approx (0.4 \text{ to } 0.5) \text{ GeV} \). The evolution of the structure function \( F_2^{p_0}(x, Q^2) \) which is implied by Eqs. (1.14) and (1.15) is shown in Figs. I.7 and I.8. We notice the characteristic shift of the structure function to low \( x \), which we earlier argued should happen in most field theories. In Fig. I.7 we have compared QCD with the predictions of the scalar and Abelian gluon field theories (1.8) and (1.9), with the value of \( g_V^q \) chosen to mimic asymptotic freedom in the range \( 2 \leq Q^2 \leq 100 \).

A crucial test of these ideas comes from the behaviour of \( R \) (1.2). In a naive parton model \( R = O(1/Q^2) \) as \( Q^2 \to \infty \), whereas in QCD it is expected that \( R = O(1/\ln Q^2) \), with a convenient approximate parametrization being

\[
R = \frac{\sigma_L}{\sigma_T} \approx \frac{1 - x}{2 \ln (Q^2/\Lambda^2)}.
\]

Data at present values of \( Q^2 \) are much larger than indicated by Eq. (1.16), which may reflect the presence of \( O(m^2/Q^2) \) pieces which have not yet died away, or may spell trouble for QCD. It is very important to test Eq. (1.16), and thereby the adequacy of quark-partons and asymptotic freedom, at larger \( Q^2 \). The formula (1.16) in fact predicts that \( R \) would be unobservably small at CHERP \( Q^2 \), although there is a connection between \( R \) and the \( p_T \) of partons and final-state hadrons (see Section 2.2). If a non-zero value were measured experimentally, it would probably require abandoning QCD or introducing boson partons (see Section 2.1.2).

If one wishes to test the general qualitative features of the scale breaking expected in field theories, experiments planned at the SPS or FNAL may suffice. But to test quantitatively different field theories of the strong interactions, experiments at very large \( Q^2 \) seem to be necessary, as seen in Figs. I.7 and I.8. As another example, Table I.1 compares different forms of scale breaking which are fixed to behave in the same way between \( Q^2 = 2 \) and 30 GeV\(^2\). It is clear that only a long lever arm like that afforded by

<table>
<thead>
<tr>
<th>( Q^2 ) (GeV(^2))</th>
<th>2</th>
<th>30</th>
<th>100</th>
<th>2000</th>
<th>10(^4)</th>
</tr>
</thead>
<tbody>
<tr>
<td>( (\ln 8/\ln 4Q^2)^{0.43} )</td>
<td>1</td>
<td>0.7</td>
<td>0.63</td>
<td>0.53</td>
<td>0.50</td>
</tr>
<tr>
<td>( (2Q^2)^{0.13} )</td>
<td>1</td>
<td>0.7</td>
<td>0.60</td>
<td>0.41</td>
<td>0.33</td>
</tr>
<tr>
<td>( (143 + Q^2)^2 )</td>
<td>1</td>
<td>0.7</td>
<td>0.35</td>
<td>4.5 \times 10^{-3}</td>
<td>2 \times 10^{-8}</td>
</tr>
</tbody>
</table>

a) The logarithmic function is that controlling the area under the valence quark distributions in asymptotic freedom with \( \Lambda^2 = 0.25 \). The other functions are chosen to behave in the same way between \( Q^2 = 2 \) GeV\(^2\) and 30 GeV\(^2\).
CHEEP can discriminate between (for example) the power-law behaviour characteristic of theories which are not [(1.8) and (1.9)] or are [(1.10) to (1.13)] asymptotically free. As a specific example, it should be possible to confirm or refute definitively the unambiguous predictions (1.10) to (1.13) by which QCD stands or falls.

2.1.2 Structure for quarks?

It may well be that the field theoretical suggestions of the previous section are completely irrelevant. Historically, constituents which seemed to be fundamental when viewed on one distance scale have turned out to be made up of smaller constituents when viewed on a smaller distance scale. There is no reason known to us why some new level of structure should appear between \(10^{-15}\) and \(10^{-16}\) cm, but the possibility exists and should therefore be pursued. Some of the proposals for quark structure will now be mentioned.

It has been suggested\(^{25}\) that quarks may have a non-trivial form factor, which might, for example, be parametrized as if it had a ('gluon') pole in it: in this case

\[
F_1(x,Q^2) = F_1(x,Q_0^2) \left( \frac{m^2 + Q_0^2}{m^2 + Q^2} \right)^2.
\]  

(I.17)

This would predict universal scale breaking at different values of \(x\), which is not the trend of the present data\(^{12}\), but might be masked by threshold effects. In Table I.1 we have given\(^2\) a form of the type (I.17) which mimics logarithmic scale breaking in the range \(2 < Q^2 < 30\) GeV\(^2\). Any such effect would be visible at CHEEP if the parameter \(m\) were \(\leq 100\) GeV. Since it would show up in the one-photon exchange diagram (I.1), it could be distinguished from any possible \(W^2\) and \(Z^0\) propagator effects.

Another possibility is that the colour degree of freedom of quarks may be excited at sufficiently small distances\(^{26}\), and the "presently observed" fractionally charged quarks resolved into integrally charged quarks of different colours:

\[
\begin{align*}
\{ & u(Q = \frac{2}{3}) \to u_R(Q = 1), \quad u_Q(Q = 1), \quad u_B(Q = 0) \} \\
& d(Q = -\frac{1}{3}) \to d_R(Q = 0), \quad d_Q(Q = 0), \quad d_B(Q = -1) \\
\end{align*}
\]  

(I.18)

The electroproduction cross-section would then rise dramatically above colour threshold, with the structure functions corresponding to individual flavours rising as\(^{27}\)

\[
\begin{align*}
\left( \frac{2}{3} \right)^2 u(x) & \to \frac{1}{3} \left( 1^2 + 1^2 + 0^2 \right) u(x) \\
\left( -\frac{1}{3} \right)^2 d(x) & \to \frac{1}{3} \left( 0^2 + 0^2 + 1^2 \right) d(x) \\
\end{align*}
\]

\[
\begin{align*}
\left\{ \begin{array}{c}
\frac{4}{3} u(x) + \frac{2}{3} u(x) \\
\frac{1}{3} d(x) + \frac{1}{3} d(x)
\end{array} \right\}
\]  

\[
\begin{align*}
\text{i.e.,} & \quad \left\{ \begin{array}{c}
u(x) \\
d(x)
\end{array} \right\}
\end{align*}
\]  

(I.19)
so that in the valence quark approximation \( u(x) = 2d(x) \), the e + p \( \rightarrow e' + X \) cross-section would rise by \( \sim 67\% \) as the colour threshold was crossed. An intriguing possibility in such a "colour liberation" model is that the photon may strike charged gluons [for example \( Q(\rho G) = 1 \)], resulting in a contribution to \( \sigma_L \) which did not vanish as \( Q^2 \rightarrow \infty \). It is clearly fundamental to know whether colour is an exact or approximate symmetry of quarks, and QED can study this question down to distances \( O(10^{-3}) \) of the nucleon radius, where weak and electromagnetic interactions are of comparable strength.

By analogy with the decomposition of previous "fundamental" constituents such as atoms and nuclei into smaller subunits, it may be expected that quarks will also turn out to be composite. The first signature for such an effect may be the excitation of quarks into higher states \( q^* \). Unfortunately, such excited quarks would probably not show up as resonance peaks in the invariant mass of the final-state hadron system, because the "Fermi motion" of the quarks inside the nucleon would be too high as to blur them out. On the other hand, the \( q^* \) might show up as a jet with a high mass (10 GeV < \( m_{q^*} \) < 100 GeV?) and width (the relationships between the constituents of the quarks are presumably very strong because they are tightly bound). Such a \( q^* \) might produce a dramatic final-state signature when it decayed back into the lowest-mass quark, emitting perhaps gluons or many pions. Eventually one might expect to decompose the quarks into the next layer of constituents, just as previous resonances (e.g. N's) have eventually been shown (in deep inelastic scattering) to be made of constituents (quarks). Each "subquark" would carry only a fraction of the longitudinal momentum \( x \) of its parent quark, so that the support of the structure function would be concentrated at small \( x \). Whether the structure function rose or fell would depend on the charges of the constituents:

\[
\begin{align*}
\{ Q = \frac{2}{9} \} & \Rightarrow Q^2 = \frac{4}{9} \Rightarrow \bar{Q}^2 = \frac{4}{27} \\
\{ Q = \frac{4}{9} \} + 2u_{s} \{ Q = -\frac{1}{3} \} & \Rightarrow Q^2 = \frac{4}{9} \Rightarrow \bar{Q}^2 = 2
\end{align*}
\]

At present, possibilities like these are totally speculative, but we will only be able to claim that we understand the unification of weak and electromagnetic interactions at the expected scale of \( O(100) \) GeV if we are also able to resolve quark structure at the same scale.

2.2 The final state in deep inelastic collisions

Looking at the deep inelastic final states will cast light on two important problems for the quark model of hadrons: what is the distribution of quark-partons inside the nucleon, and how a parton recombines into hadrons after it has been given a high momentum kick. We will discuss possibilities for these processes first in the context of the naive quark parton model, and then in the light of recent developments based on field theories such as QCD.

2.2.1 The naive parton picture

In this approach, quark-partons are required to have finite momenta transverse to the proton beam direction. Then, when struck by a virtual \( \gamma \) (or \( W^\pm \) or \( Z^0 \)), they are kicked
in the direction of the momentum transferred from the leptons, and finally convert into hadrons with finite momenta transverse to their new momentum axis. Meanwhile, the parton constituents of the nucleon which were not struck still proceed in the proton beam direction, and eventually recombine into hadrons with finite transverse momenta. Thus all the final-state hadrons should occupy a finite width "cigar" in the Feynman plot aligned parallel to the directions of the proton and virtual photon in their centre-of-mass frame, as illustrated in Fig. I.9. Bjorken\(^{24}\) has distinguished five regions of longitudinal phase space which are illustrated in Fig. I.10a. Within a finite rapidity width \(dy = O(2)\) of the kinematic boundary in the proton direction, \(y = Y_1\), should lie hadrons which have fragmented from the high-momentum partons left in the proton after the deep inelastic collision. If the struck parton had very low \(x\), then as seen from Figs. I.10a and I.10b, these hadron fragments should be identical with those seen in high-energy \(p+p\) collisions, as at the Intersecting Storage Rings (ISR), with a proton fragmentation region and a plateau like that observed in the central region of hadron-hadron collisions:

\[
\frac{1}{d\sigma(\gamma^*p)/dx} \frac{d^2\sigma}{dxdy} (\gamma^* + p \to H + X) = \frac{1}{\sigma_{\text{tot}}(pp)} \frac{d\sigma}{dy} (p + p \to H + X) \quad (I.21)
\]

for \(|Y_1 - y| < \ln(1/x)\). In the neighbourhood of rapidity space whence the struck parton was removed \(|Y_1 - y| = \ln(1/x) \pm O(2)\), there should be a "hole fragmentation" region reflecting the behaviour of a hadron which has had a pole punched in its wave function. If the struck parton does not have very low \(x\), the hole and proton fragmentation regions will run together to give a distribution which should be identical\(^{25}\) with that found in the hadronic final state in a Drell-Yan \(p + p \to \gamma^*\), \(W\), or \(Z + X\) reaction (see Fig. I.10c):

\[
\frac{1}{d\sigma(\gamma^*p)/dx} \frac{d^2\sigma}{dxdy} (\gamma^* + p \to H + X) = \frac{1}{\sigma[p + p \to \ell\bar{\ell}(\tau, z) + H + X]} \frac{d\sigma}{dy} [p + p \to \ell\bar{\ell}(\tau, z) + H + X] .
\]

The quantities \(\tau \equiv m^2(\ell\bar{\ell})/s\) and \(z \equiv 2(E_\ell + E_{\bar{\ell}})/\sqrt{s}\) must be chosen such that

\[
x(z-x) = \tau ,
\]

which ensures that the parton removed from the proton by the Drell-Yan mechanism has the same value of \(x\) as the struck parton in the electroproduction reaction.

At the opposite end of rapidity space: \(|Y_2 - y| = O(2)\) in Fig. I.10a there should be a finite length quark-parton fragmentation region identical with that seen in lower-energy \(e^-p\), \(e^-p\), and \(\nu^-p\) collisions, and supposedly\(^{26}\) the same as in \(e^+e^-\) collisions\(^{27}\) (Fig. I.10d):

\[
\frac{1}{d\sigma(\gamma^*p)/dx} \frac{d^2\sigma}{dxdy} (\gamma^* + p \to H + X) = \frac{1}{\sigma(e^+e^- \to \text{hadrons})} \frac{1}{d\sigma/ dy} (e^+e^- \to H + X) \quad (I.24)
\]

\(^{24}\) We ignore for the time being the complications engendered by having several different types of quark with different quantum numbers.
if we again ignore complications due to different quark quantum numbers. In between the hole and quark-parton fragmentation regions there is supposed to be a second "current" plateau, again identical to that found in $e^+e^-$ collisions, so that the identity (I.24) should hold for all $y > Y_1 + \ln(1/x)$.

To see all the five kinematic regions distinguished above would probably need a total rapidity length $O(10)$ with $Q^2 \gtrsim 100$ GeV$^2$ in order to see the current plateau, which just starts to manifest itself at $Q^2 = 50$ GeV$^2$ in $e^+e^-$ annihilation. Reference to Fig. I.2 shows that thousands of events per day may be expected to meet these kinematical criteria. When these regions are distinguishable, it should be possible to make interesting checks of the basic parton picture by studying quantum number correlations. For example, if an $s$ quark is struck by the virtual photon, the strange particle in its fragmentation may have its strangeness compensated by a strange antiparticle coming in the hole fragmentation region from the $\bar{s}$ quark left behind in the photon wave function. Their rapidity separation would be

$$\Delta y = Y_2 - \left[-Y_1 + \ln(1/x)\right] = \ln Q^2,$$  \hspace{1cm} (I.25)

whereas in normal hadron-hadron collisions, or in the usual final-state recombination in electroproduction collisions, we would expect to find $\Delta y = O(2)$. We may therefore expect a distribution for $\Delta y$ like that shown in Fig. I.11, for other quark quantum numbers such as charm as well as for strangeness. A similar effect may be visible with electric charge. Since the virtual photon most often strikes a $u$ quark in the proton, and the fragmentation $u \rightarrow \pi^+$ is favoured over $u \rightarrow \pi^-$, it is expected that there would be a net positive charge in the parton fragmentation region. Correspondingly there should be a net negative charge in the hole fragmentation region. The correlation would of course be reversed for the case where a $d$ quark is struck by the virtual photon. Therefore if we plot the quantity

$$\frac{d}{dy} \left[ Q_\perp \text{sign}(Q_\perp) \right],$$

where $Q_\perp$ is the net charge in the parton fragmentation region, then we might expect a structure like that in Fig. I.11, indicating a non-local compensation of charge not found in hadron-hadron collisions.

For very large values of $W^2 \equiv (q+p)^2, 1/x,$ and $Q^2$, we expect the development of the hadronic and current plateaux $^{28)}$ to imply $^{31)}$ that the multiplicities will be related by

$$\langle n(W^2) \rangle_{ep} = \langle n[(1-x)W^2] \rangle_{pp} + \langle n(Q^2) \rangle_{e^+e^-} + \text{finite corrections}, \hspace{1cm} (1.26)$$

where the corrections come from the various fragmentation regions. There is some dispute as to whether the plateaux heights $(1/a)(d\sigma/dy)$ should be identical in the different reactions. Simple-minded universality $^{32)}$

$$\frac{\text{hadronic height}}{\text{current height}} = 1$$  \hspace{1cm} (1.27a)
may be a good approximation to present-day data, although other predictions exist. For example, in dual models\(^1\) and some quark models\(^1\) we expect two overlapping jets of hadrons in the hadronic plateau region (see Fig. 1.12), so that

\[
\frac{\text{hadronic height}}{\text{current height}} = 2. \quad (1.27b)
\]

Another model\(^2\) focuses on the amount of colour separation taking place in the process: \(\gamma \rightarrow \overline{\gamma}\) in the case of \(e^+e^- \rightarrow q\overline{q}\) or electron production in the current plateau, possibly \(g \rightarrow g\) in hadron-hadron collisions because the Pomeron may correspond to octet gluon exchange. If the plateau height is proportional to the colour field density, we expect\(^3\)

\[
\frac{\text{hadronic height}}{\text{current height}} = 2^{1/4}. \quad (1.27c)
\]

Present data are at energies too low to discriminate decisively between the different predictions (1.27): if either of the ideas (1.27b) or (1.27c) turn out to be correct, then the hole fragmentation region should be a very dramatic transition region.

Some comments on the final-state kinematics may be in order here. In an average CHEP e-p collision, the x of the struck parton will probably be \(\mathcal{O}(1/10)\), in which case the collision will resemble an electron-quark collision at \(\mathcal{O}(50)\) GeV in the centre of mass. At large \(Q^2\), the parton and its current jet will emerge at a large angle relative to the beam axes. On the other hand, the proton fragmentation, hadronic plateau, and hole fragmentation regions would continue close to the direction of the incoming proton beam, but with finite \(p_T\). Thus, although much of the interesting physics lies in the current region, the simple parton model\(^4\) discussed here suggests that there may be interesting and important dynamics to be uncovered in other, less experimentally accessible, regions of final-state phase space.

### 2.8.2 Possible modifications to the naive parton picture

To date there is very little evidence for the elegant picture outlined above, except that there is evidence for finite momenta of final-state hadrons transverse to the momentum transferred from the lepton system in deep inelastic collisions, and indications exist\(^5\) of the universality of quark-parton fragmentation in e-p, \(\nu\overline{\nu}-p, \nu-p,\) and \(e^+e^-\) collisions. Confirmation of these and the other predictions discussed above would be a great success for simple-minded quark-parton ideas\(^6\). However, just as experiment and theory strongly suggest that there are deviations from the simple parton model in the deep inelastic structure functions, so there may also be different effects in the hadronic final state. Several possibilities come to mind in the context of field theory.

It is well known that in most field theory models the \(p_T\) of the quanta in a deep inelastic collision are not finite as \(Q^2 \rightarrow \infty\). Generally, the momenta transverse to the scattering plane would behave as

\[
\langle p_T^2 \rangle = \mathcal{O}(\ln^n Q)
\]
in any finite order of perturbation theory. In the parton language of the previous section this might either reflect the growth of the $p_T$ of the quark-parton, or the growth of $p_T$ in the parton fragmentation region, or both effects. In a weak binding approximation the initial parton $p_T$ is connected with $R = \sigma_L/\sigma_T$:

$$\frac{\sigma_L}{\sigma_T} \approx \frac{4(p_T^2)}{Q^2} \quad (I.28)$$

which, combined with the result (I.16) sanctioned by the renormalization group, suggests that in an asymptotically free gauge theory$^{23)}$

$$\langle p_T^2 \rangle \approx \frac{(1-x)Q^2}{8 \ln (Q^2/\Lambda^2)} \quad (I.29)$$

In fact, asymptotically free theories are not completely free, so that the approximation (I.28) may be invalid. Nevertheless, Eq. (I.29) probably gives a reasonable order of magnitude for $\langle p_T^2 \rangle$: it would suggest $\langle p_T \rangle = O(5 \text{ to } 10) \text{ GeV at } Q^2 = O(10^4) \text{ GeV}^2$.

Another effect$^{39)}$ which might be anticipated in field theory is that, just as the deep inelastic structure function $F_2(x,Q^2)$ should fall in towards $x=0$ as $Q^2 \to \infty$, so also the quark to hadron fragmentation function $D(z,Q^2)$: $z \equiv E_{\text{hadron}}/E_{\text{quark}}$ might also move in towards $z=0$ as $Q^2 \to \infty$ as indicated in Fig. I.13. This behaviour would be expected in a picture like that of Gribov and Lipatov$^{39)}$ where some sort of continuation between electron production and $e^+e^-$ annihilation cross-sections is possible$^{39)}$.

There is at present a widespread theoretical effort to provide a firm theoretical basis for these considerations, which may be useful to stretch the mind, even though the conclusions are tentative at the moment. The approach considers cross-sections for the production of a certain number of quarks and gluons as "jet cross-sections" corresponding to a certain amount of hadronic energy being deposited in a specified range of solid angle$^{39,39,37-39)}$. It is believed that suitably defined jet cross-sections are infrared finite$^{39)}$ and obey renormalization group equations in field theory$^{41)}$, in which case they may easily be calculated using the strong interaction coupling constant $g$ discussed in Section 2.1. For example, the prediction for three-jet events in electroproduction is that one jet goes in the direction of the incoming $\gamma^*$, while the distribution of the other two jets is given by$^{39,39,42)}$:

$$\frac{d^3\sigma}{dx dy p d\phi} = \sum_q q(x,Q^2) \left\{ \frac{d^3\sigma}{dx dy p d\phi} (\gamma^* + q + g + q) \\
+ g(x,Q^2) \sum_q \frac{d^3\sigma}{dx dy p d\phi} (\gamma^* + g + q + \bar{q}) \right\} \quad (I.30)$$
In Eq. (I.30),
\[ x_T = \frac{2p_T}{W} : W = (q + p)^2 \]  
(1.31)

and \( \phi \) are the scaled transverse momentum and azimuthal angle, respectively, of the final-state jets relative to the direction of the incoming \( \gamma^* \). In an asymptotically free gauge theory the cross-sections for \( \gamma^* + q \rightarrow g + q \) and \( \gamma^* + g \rightarrow q + \bar{q} \) would be calculated in lowest-order perturbation theory as indicated in Fig. (I.14); the \( q(x, Q^2) \) are the asymptotically free quark distributions parametrized in Eqs. (I.14) and (I.15), and \( g(x, Q^2) \) is the analogous gluon distribution.

Equation (I.30) determines a scaling law for high \( p_T \) cross-sections; several remarks should be made about this prediction:

i) It is, as yet, on a less secure footing than the QCD results for the moments of the structure functions discussed in Section 2.1.

ii) Similar predictions would exist in other field theories which are not asymptotically free. They would generally predict larger three-jet cross-sections than the fraction \( O(\alpha_s/\pi) = O(10\%) \) expected from Eq. (I.30), just because \( g \rightarrow g^* \neq 0 \) in such theories.

iii) Predictions analogous to (I.30) could be made for four or more jet cross-sections.

In an asymptotically free gauge theory, such a "jet perturbation expansion" is probably reasonably convergent.

iv) Note that since the large \( p_T \) cross-section in (I.30) is obtained by summing the two gauge-variant diagrams in Fig. (I.14), it is difficult to separate meaningfully the contributions from initial partons with large \( p_T \) and final-state fragmentation with large \( p_T \). This means that the formula (I.29) \(^{23}\) should perhaps be taken with a pinch of salt.

v) While Eq. (I.30) refers to the cross-section for producing a certain value of \( \sum |p_T| \) summed over all particles in the event, analogous calculations can be made for the one-particle inclusive cross-sections at large \( p_T \) if we assume forms for the parton to hadron fragmentation functions. A calculation of this type \(^{24} \) is shown in Fig. I.15. The prediction is much larger than the high \( p_T \) cross-section found at the ISR in hadron-hadron collisions.

The above ideas based on field theory suggest considerable richness in the hadronic final state beyond that already suggested by the naive parton model of Section 2.2.1. Whatever picture of the hadronic final state turns out to be correct, it is clear that studies of it will be an important tool for revealing how partons are contained in hadrons, and how partons eventually change back into hadrons. At the moment we have much less experimental and theoretical understanding of these questions than of the deep inelastic structure functions themselves -- but these questions are very important because they cut straight to the heart of the problem of quark confinement.

2.3 New quark production

New quarks may be pair-produced in deep inelastic electroproduction. It seems likely that as \( Q^2 \rightarrow \infty \) the production rates of new massive quarks should increase relative to that of low-mass quarks. This expectation is common to many approaches -- field-theoretic and
others such as vector meson dominance--but there is very little experimental indication of its validity\(^{(11)}\). The basic reasoning is that as \(Q^2 \to \infty\), one is probing smaller space-time regions within hadronic matter, and the correlation with the global quantum numbers of the hadron target should vanish as one probes more and more locally. A specific field-theoretic mechanism\(^{(16)}\) is depicted in Fig. 1.5, where new quark production is visualized as pair production in the "gluon" field of the nucleon target, analogous to the Bethe-Heitler pair production process in QED. At low \(Q^2\), a "large" low resolving power photon fails to dissociate a "gluon", whereas at higher \(Q^2\) the "gluon" can be resolved into a quark-antiquark pair. For \(Q^2 >> m_{\text{quark}}^2\), we might expect this process to be independent of \(m_{\text{quark}}\), and hence SU(F) symmetric at large momentum transfers, where \(F\) is the number of different quark flavours.

Such processes may be expected in any field theory. Just as scaling violations generally are larger in theories which are not asymptotically free, we expect the pair production of new quarks to be minimized in an asymptotically free gauge theory (QCD). Table 1.2 shows the result of calculations in QCD\(^{(12)}\) for the production at \(x = 0.05\) of different types of high-mass quarks at different values of \(Q^2\). The trend towards SU(F) symmetry is apparent. Similar results are found in a generalized vector meson dominance approach\(^{(14)}\), where the relative production rates of new quarks are somewhat larger than the QCD estimates.

**Table 1.2**

<table>
<thead>
<tr>
<th>Characteristics of new quark a)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(Q^2 \text{ (GeV}^2))</td>
</tr>
<tr>
<td>50</td>
</tr>
<tr>
<td>100</td>
</tr>
<tr>
<td>500</td>
</tr>
<tr>
<td>1000</td>
</tr>
<tr>
<td>(\infty)</td>
</tr>
</tbody>
</table>

a) The above are percentages of the total electromagnetic cross-section calculated\(^{(21)}\) for \(x = 0.05\). Generalized vector dominance gives somewhat larger percentages of new heavy quark production.

The kinematics of these events should be quite favourable, with one of the heavy quark mesons emerging in the current jet, and hence at a large angle. It has been widely suggested\(^{5,6}\) that the heavy quark meson \(H\) would generally carry a large fraction

\[
(z) \sim 1 - \frac{1}{m_Q \text{(GeV)}}, \quad z \equiv \frac{E_H}{E_Q}
\]

(I.32)

of the total energy of the heavy quark jet, expected to be \(\mathcal{O}(10 \text{ to } 50)\) GeV, with the rest of the energy carried by slow-moving light hadrons, as indicated in Fig. 1.16a. A possible
signature for such an event is the production of a prompt lepton in a semileptonic decay of \( H \), which is expected to occur with a branching ratio \( O(10 \text{ to } 20\% \). Such prompt decay leptons are expected to carry \( O(1/3) \) of the heavy quark energy (see Fig. 1.6b), and to have large momenta transverse to the current jet which are \( O(m_H/3) \) (see Fig. 1.17) \(^{47}\). Figure 1.6b also shows the longitudinal momentum distribution expected\(^{46}\) for light quark hadrons produced in the semileptonic decay. In a non-leptonic decay, the dominant decay modes of the heavy meson \( H \) are expected to proceed via \( Q \to q\bar{q}Q' \), where the \( q \) are light quarks. This process should\(^{40}\) yield high multiplicity final states with all hadron longitudinal momenta \( \leq E_Q/3 \) as seen in Fig. 1.16c. These final states would also have large total hadron momenta transverse to the current jet axis:

\[
\sum_{\text{hadrons } i} |p_i^\perp| = O(m_Q) . \tag{1.33}
\]

A point worth noting is that heavy quarks may well tend to cascade in their decays to lighter quarks: \( t \to bff, b \to cff, c \to sff = q\bar{q} \) or \( \bar{q} \bar{u} \), for example, which could provide multilepton final states.

The above remarks suggest several handles for identifying and analysing final states likely to contain new heavy quark hadrons. Also, Table I.2 shows that the signal-to-background ratios for new quark production may be only a factor (2 to 4) worse than in \( e^+e^- \) annihilation above threshold, while the kinematics are similarly favourable. Electroproduction is not expected to provide a kinematic region analogous to the (4.0 to 4.4 GeV) region of enhanced charm production in \( e^+e^- \) annihilation, but should provide a good laboratory for studying the production and decays of new heavy quarks.

### 2.4 Photoproduction

As discussed elsewhere in this report, there should be more than \( 10^6 \) events per day at CHEEP, of the form \( e + p \to e' + X \) with \( Q^2 \leq m_e^2 \), which can be used to study photoproduction at equivalent beam energies up to 20 TeV. One of the most striking features of photoproduction studies to date is the extent to which a real photon interacts as if it were a normal hadronic vector meson with its coupling reduced by \( \sim 1/250 = O(\alpha/\pi)^{48} \). If this still happens at beam energies two orders of magnitude higher than those obtained from present-day accelerators, one may identify the following areas of photoproduction investigations:

- the study of "hadron"-hadron collisions at centre-of-mass energies beyond those of the ISR;
- the search for unhadronic, distinctly photonic effects at \( Q^2 = 0 \);
- the use of the specific quantum numbers of the photon, \( J^{PC} = 1^{--} \), to produce diffractively new hadronic states, and the similar use of the Primakoff process to study new \( C = +1 \) states.

As far as the first topic is concerned, there are two interesting features of photon-hadron scattering which make it especially interesting. The first is that the beam has a non-zero, variable, and specifiable polarization. The second is that the Froissart bound does not apply to \( \sigma_{\text{tot}}(\gamma p) \), which adds to the interest of the energy-dependence of the expected rise in \( \sigma_{\text{tot}}(\gamma p) \), and the comparison with the increases in hadron-hadron total...
cross-sections. One mechanism which may contribute to an abnormal rise in $\sigma_{\text{tot}}(\gamma p)$ is the production of new quark flavours. The total cross-section for charm pair photoproduction is expected to be $O(1)$ $\mu$b, or about 1% of the total cross-section, a much larger fraction than in hadron-hadron collisions. The cross-sections for new heavier quark flavours are presumably smaller, but may still amount to some sizeable fraction of a microbarn.

It has often been pointed out that one of the most fundamental processes to be studied with a photon beam is Compton scattering. It may reveal properties unique to the photon; for example, a fixed pole which is disallowed in purely hadronic reactions, but is expected to dominate Compton scattering at high energies and large values of $t$. The parton model of Fig. 1.18 gives an almost real amplitude, and hence a large difference between the cross-sections for photons with their polarization vectors either in or normal to the scattering plane.

A strong analogue to Compton scattering may also exist, namely the continuation to $Q^2 = 0$ of the "gluon" bremsstrahlung process illustrated in Fig. 1.14. In the case of a non-Abelian vector "gluon", Fritzsch and Minkowski estimate

$$\frac{d^2\sigma}{dt d\gamma^2} = \frac{16\alpha_s}{\pi^2 t} \cdot \frac{\alpha_s}{s} \left( \frac{u}{s} + \frac{s}{u} \right) \frac{F^2_{\gamma p}(x)}{x},$$

where, if we define $k$, $p$, and $k'$ as the photon, proton, and outgoing gluon momenta, respectively:

$$v' \equiv p \cdot k', \quad s \equiv (k+p)^2, \quad t \equiv (k-k')^2, \quad u \equiv (p-k')^2, \quad x \equiv \frac{-t}{2(p \cdot k - u')},$$

and $\alpha_s = g^2/4\pi$ is a strong coupling $O(0.1$ to 1). They suggest that $O(10^{-4}$ to $10^{-5})$ of the total photon-nucleon cross-section for $p_{\text{lab}} \approx 200$ GeV may contain events with recognizable gluon and quark jets emerging at large centre-of-mass angles. This "gluon bremsstrahlung" model is not the only one where large $p_T$ cross-sections are expected to be different in hadron- and photon-induced processes; such effects are also expected in constituent interchange models.

Recent experience suggests that photon-hadron collisions differ from hadron-hadron collisions in having a larger fraction of new heavy particle production. The reason for this belief is that the photon couples to even the heaviest quark through its electromagnetic charge. Thus the lowest-lying vector meson bound states of new heavy quarks can be diffractively produced. Table 1.3 lists the production rates of the known light vector mesons, and guesses for heavier mesons. Generally, vector meson dominance suggests that

$$\frac{d\sigma}{dt} (\gamma + p \to V + p) \approx \left( \frac{m_V^2}{2} \right) \frac{d\sigma}{dt} (V + p \to V + p).$$

If high-mass vector mesons are produced diffractively as in Fig. 1.19, and have similar slopes for their diffraction peaks, Eq. (1.36) implies that

$$\frac{\sigma(\gamma^* p \to V^* p)}{\sigma(\gamma^* p \to J/\psi p)} = \frac{\gamma_J^2}{\gamma_V^2} \left[ \frac{\sigma(V^* p \to X)}{\sigma(J/\psi p \to X)} \right]^2.$$
Table I.3
Photoproduction of vector mesons *)

<table>
<thead>
<tr>
<th>Reaction</th>
<th>$\sigma_{\text{tot}}$</th>
<th>Events/day tagged</th>
<th>Events/day no tag</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\gamma + p \to \rho^0 + p$</td>
<td>13 $\mu$b</td>
<td>$5.6 \times 10^5$</td>
<td>$11 \times 10^6$</td>
</tr>
<tr>
<td>$\gamma + p \to \omega + p$</td>
<td>1.4 $\mu$b</td>
<td>$6 \times 10^8$</td>
<td>$1.2 \times 10^8$</td>
</tr>
<tr>
<td>$\gamma + p \to \phi + p$</td>
<td>0.5 $\mu$b</td>
<td>$2.2 \times 10^9$</td>
<td>$4.4 \times 10^9$</td>
</tr>
<tr>
<td>$\gamma + p \to J/\psi + p$</td>
<td>30 nb</td>
<td>$1.3 \times 10^3$</td>
<td>$2.6 \times 10^4$</td>
</tr>
<tr>
<td>$\gamma + p \to T + p$</td>
<td>$\mathcal{O}(50)$ pb?</td>
<td>$\sim 2$</td>
<td>$\sim 40$</td>
</tr>
<tr>
<td>$\gamma + p \to V(20 \text{ GeV}) + p$</td>
<td>$\mathcal{O}(1)$ pb?</td>
<td>$\sim 0.04$</td>
<td>$\sim 1$</td>
</tr>
</tbody>
</table>

*) Rates above the dashed line are known, rates below are guesses.

If

$$\sigma(V + p \to X) = \frac{1}{m_V^2}$$

(1.38)

and

$$\frac{\gamma^2_{J/\psi}}{\gamma^2_V} \approx \left(\frac{\mathcal{O}_{\text{new quark}}}{\mathcal{O}_{\text{charm}}}\right)^2 \frac{m_{J/\psi}}{m_V}$$

(1.39)

then

$$\frac{\sigma[\gamma + p \to T(9.5) + p]}{\sigma(\gamma + p \to J/\psi + p)} \approx \left(1 \text{ or } \frac{1}{4}\right)\left(\frac{1}{300}\right)$$

(1.40)

depending on the charge ($\gamma^0$ or $-\gamma^0$) of the quarks in the $T$. One might therefore guess

$$\sigma(\gamma + p \to T + p) = \mathcal{O}(50) \text{ pb}$$

(1.41)

to within an order of magnitude. A hypothetical bound state of a quark of mass 10 GeV might have a cross-section one or two orders of magnitude less than (1.41): hence the guesses in the last two lines of Table I.3. Even vector mesons with masses $\mathcal{O}(20)$ GeV, corresponding to rates $\mathcal{O}(1)$ per day, might be detectable because of their distinctive dilepton or multi-hadron jet decay signals.

New even-charge conjugation states with $J \neq 1$ can be produced by the Primakoff effect shown in Fig. I.20. The cross-section is

$$\frac{d\sigma}{dt} (\gamma + p \to X + p) = \frac{8\pi\alpha}{m_X^3} \Gamma(X \to \gamma\gamma) \left(\frac{t-t_{\text{min}}}{t^2}\right) \left[F_p(t)\right]^2,$$

(1.42)

where $F_p(t)$ is the proton electromagnetic form factor and

$$t_{\text{min}} = \frac{m_X^2 m_p^2}{s}.$$

(1.43)
If $t_{\text{max}}^{\text{eff}}$ is the effective momentum transfer cut-off due to the proton form factor, Eq. (I.42) yields a total cross-section (for $t_{\text{min}} < t_{\text{max}}^{\text{eff}} < s$):

$$\sigma(\gamma*p\rightarrow X*p) = \frac{8\pi}{m_X^2} \Gamma(X\rightarrow \gamma\gamma) \ln \frac{t_{\text{max}}^{\text{eff}}}{t_{\text{min}}}$$  \hspace{1cm} (I.44)

which increases logarithmically with energy. The rates for some particles with large $\gamma\gamma$ decay widths are listed in Table I.4. One might hope to look for the various $C = +1$ charmonium states, as well as the pseudoscalar partner of the $\Upsilon(9.5)$.

### Table I.4

<table>
<thead>
<tr>
<th>Particle</th>
<th>$\Gamma_{\gamma\gamma}$ (MeV)</th>
<th>$\sigma$ (nb)</th>
<th>Events/day</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\eta^*$</td>
<td>$8 \times 10^{-6}$</td>
<td>1.4</td>
<td>1100</td>
</tr>
<tr>
<td>$\eta$</td>
<td>$3.8 \times 10^{-6}$</td>
<td>0.99</td>
<td>560</td>
</tr>
<tr>
<td>$X^0$</td>
<td>$2 \times 10^{-2}$</td>
<td>9.8</td>
<td>4710</td>
</tr>
<tr>
<td>$f$</td>
<td>$\sim 4 \times 10^{-3}$</td>
<td>0.84</td>
<td>370</td>
</tr>
<tr>
<td>$n_C$</td>
<td>$\sim 20 \times 10^{-3}$</td>
<td>0.59</td>
<td>90</td>
</tr>
</tbody>
</table>

3. **WEAK INTERACTION STUDIES**

If weak interaction cross-sections continue to rise linearly with laboratory energy like they do at present, then they will become comparable with electromagnetic cross-sections at centre-of-mass energies $O(50$ to $100)$ GeV:

$$\frac{G_F}{\sqrt{2}} = O(e^2/Q^2) \quad \text{for} \quad Q^2 = O(10^3$ to $10^4)$ \hspace{1cm} (I.45)

It is now considered likely that the weak and electromagnetic interactions cannot be considered independently, but must instead be unified into a combined theory. Such theories generally incorporate intermediate vector bosons which damp the traditional four-fermion interaction amplitude

$$\frac{G_F}{\sqrt{2}} \rightarrow \frac{g^2}{8(m_W^2-Q^2)}.$$ \hspace{1cm} (I.46)

The vector bosons typically have masses $O(50$-$100)$ GeV, and are generally put into a gauge theoretical model\(^6\). Such a unification has no direct experimental support, and alternatives are readily conceivable. In principle, weak cross-sections could continue to rise until they violate unitarity, which happens at centre-of-mass energies $O(500$ to $1000)$ GeV. Or perhaps weak cross-sections neither rise nor fall above $E_{\text{cm}} = 100$ GeV, but are essentially constant.

It is therefore very interesting to look at weak interaction effects in a momentum transfer range where they are comparable with electromagnetic ones. We may hope to
- see whether weak and electromagnetic interactions can indeed by described by some unified theory\(^6\);
- see whether there are deviations from four-fermion cross-sections which might reflect intermediate boson propagators;
- perhaps search for the intermediate bosons themselves;
- produce new leptons and quarks.

We will see that CHEEP should be able to answer the first two of these questions. It will probably be rather difficult to actually detect \(W^\pm\) or \(Z^0\) bosons with this machine, but this possibility should not be excluded. We can reasonably expect to produce new quarks, or leptons coupled to the electron, which have masses up to 50 or 100 GeV.

### 3.1 Neutral current effects

We now go beyond the one-photon contribution to the cross-section for \(e + p \rightarrow e' + X\) to include the exchange of neutral gauge bosons \(Z_i^\pm\), as shown in Fig. 1.21. One or more of these may exist and interfere with photon exchange. In general\(^{51}\), there will be violation of parity and a charge asymmetry \([\sigma(e^- + p \rightarrow e^- + X) \neq \sigma(e^+ + p \rightarrow e^+ + X)\] which may be \(O(1)\) at \(Q^2 = O(10^2)\). In discussing these effects it is convenient to use the reference cross-section

\[
dl{\frac{d^2\sigma_{1\gamma}}{dx dy}} = \frac{4\pi\alpha^2}{sx^2y^2} \left\{ (1-y)F_2(x,Q^2) + y^2xF_1(x,Q^2) \right\} . \tag{I.47}
\]

In a parton model, we then have

\[
dl{\frac{d^2\sigma_{1\gamma}}{dx dy}} = \frac{d^2\sigma_{1\gamma}}{dx dy} \sum_{i,j} G_{\eta}^i p_j^i p_{\eta} \left[ A_{\eta}^{ij}(x) + B_{\eta}^{ij}(x)f(y) \right] , \tag{I.48}
\]

where \(\eta\) denotes the charge and polarization state of the incoming lepton \((e^\pm_L, R)\), and the other quantities in (I.48) are defined as follows.

The \(G_{\eta}^i\) are the couplings of the \(i\)th neutral boson to the incoming lepton, with the photon corresponding to \(i = 1\), so that

\[
G_{\eta}^1 = -1 \quad \text{for} \quad \eta = e_L^-, e_R^-, e_L^+, e_R^+ . \tag{I.49}
\]

The \(P_i^\pm\) are conveniently normalized propagator factors for the \(i\) vector bosons (if they exist), so that

\[
P_1^1 = 1 , \quad P_i^i = \frac{G_{\eta}^i Q^2 m^2_{\eta}}{r^2 e^2 (m^2_{\eta} + Q^2)} \quad \text{for} \quad i > 1 . \tag{I.50}
\]

The \(A_{\eta}^{ij}(x)\) and \(B_{\eta}^{ij}(x)\) reflect the products of the couplings of \(Z_i^\pm\) and \(Z_j^\pm\) to the quark-partons:
\[ A^{ij}(x) = \frac{1}{2} \sum_q \left( G^i_{q_L} G^j_{q_L} + G^i_{q_R} G^j_{q_R} \right) \left[ q(x) + \bar{q}(x) \right] \]
\[ B^{ij}(x) = \frac{1}{2} \sum_q \left( G^i_{q_L} G^j_{q_L} - G^i_{q_R} G^j_{q_R} \right) \left[ q(x) - \bar{q}(x) \right] \]

The other factors in Eq. (I.48) are defined by
\[ \xi_\eta = +1 \quad \text{for} \quad \eta = e^+_L, \quad e^+_R, \quad \xi_\eta = -1 \quad \text{for} \quad \eta = e^-_L, \quad e^-_R, \]

and by
\[ f(y) \equiv \frac{\gamma - y^2/2}{1 - y\gamma y^2/2} \]

when we assume the Callan-Gross relation\(^{12}\) to be approximately valid. It is evident from Eqs. (I.48) and (I.51) that the unravelling of all neutral current effects will be crucially dependent on having a positron beam available, and on having the possibility of longitudinally polarized beams. The normalization of the weak currents is chosen to be

\[ J_{\gamma}^{ui} = \frac{1}{2} \sum_q \bar{q}_\mu \left[ G^i_{q_L} (1 - \gamma_5) + G^i_{q_R} (1 + \gamma_5) \right] q \]

so that the electromagnetic current
\[ J_{\gamma_{em}}^{ui} = \sum_q \bar{q}_\mu q_q \]
corresponds to
\[ G^i_{q_L} = G^i_{q_R} = Q_q \]

The formulae (I.48) to (I.56) can be used with any weak interaction model, gauge or otherwise, and any quark-parton distributions, scaling or asymptotically free or whatever. As an example of a weak interaction model we may wish to test, we may cite the simplest Weinberg-Salam\(^{12}\) model based on the group \( SU(2) \times U(1) \) with only left-handed fermion doublets. This model has only one \( Z^0 \) boson, with mass

\[ m_Z = \frac{37.4}{\sin \beta_w \cos \beta_w} \]

and couplings
\[ C^2_{\eta} = \begin{cases} (-1 + 2 \sin^2 \theta_W) & \text{for } \eta = e^-_L e^+_R, \\ 2 \sin^2 \theta_W & \text{for } \eta = e^-_R e^+_L \\ \end{cases} \]  \hspace{1cm} (I.58)

and

\[ C^2_{u_L} = C^2_{c_L} = \left( 1 - \frac{4}{3} \sin^2 \theta_W \right), \quad C^2_{u_R} = C^2_{c_R} = \left( -\frac{4}{3} \sin^2 \theta_W \right) \]  \hspace{1cm} (I.59)

for the charge \( \frac{3}{2} \) quarks,

\[ C^2_{d_L} = C^2_{s_L} = \left( -1 + \frac{2}{3} \sin^2 \theta_W \right), \quad C^2_{d_R} = C^2_{s_R} = \left( \frac{2}{3} \sin^2 \theta_W \right) \]  \hspace{1cm} (I.60)

for the charge \( \frac{1}{2} \) quarks. In order to see how effective the interference effects are as probes of the neutral currents, we have considered a couple of other weak interaction models to compare with (I.58) to (I.60). One possibility is to add to this standard SU(2) \(_L\) \( \otimes \) U(1) model an ad hoc \( \left( \begin{array}{c} L^0 \\ R \\ \end{array} \right) \) doublet\( ^{54} \) which has the effect of changing the neutral current couplings of the leptons:

\[ \text{R.H. doublet model : } C^2_{\eta} (2 \sin^2 \theta_W) \quad \text{for } \eta = e^\pm_{L,R}. \]  \hspace{1cm} (I.61)

This model is a simple way of removing the apparent conflict between the first generation of atomic physics parity violation experiments\( ^{54} \) and the simplest Weinberg-Salam model\( ^{52} \) outlined in Eqs. (I.57) to (I.60). The other model considered was the SU(2) \(_L\) \( \otimes \) SU(2) \(_R\) \( \otimes \) U(1) model of De Rújula, Georgi and Glashow\( ^{55} \) in which there are two \( Z^0 \) bosons, with masses \( 0(80, 150) \) GeV, respectively. It may well be that none of these models will still be a candidate theory of weak and electromagnetic interactions in a few years time, but they may point up some of the effects to be anticipated.

Graphs of some calculated quantities are shown in Figs I.22 to I.28. Figure I.22 shows the ratios of total inclusive cross-sections for different \( e^\pm_{L,R} \) beams to the purely electromagnetic cross-section, at \( s = 27,000 \) GeV\(^2 \) corresponding to a 25 GeV \( e^\pm \) beam and a 270 GeV p beam, and \( x = 0.25, y \) in the range 0 to 1 corresponding to \( Q^2 \) from 0 to 6750 GeV\(^2 \). We see that indeed parity violation effects \( [\sigma(e^\pm_L)/\sigma(e^\pm_R)] \) and charge asymmetries \( [\sigma(e^\pm_{L,R})/\sigma(e^\pm_{L,R})] \) are substantial over much of the \( y(Q^2) \) range. Figure I.23 shows the ratio \( \sigma(e^\pm_{L,R})/\sigma(e^\pm_R) \) in the Weinberg-Salam\( ^{52} \) model for a range of different values of \( \sin^2 \theta_W \) and the same kinematic conditions as before. Present experiments generally favour \( \sin^2 \theta_W = 0.25 \) to 0.55, and these extremes seem to be readily distinguishable. Figure I.24 shows the ratio \( \sigma(e^\pm_L)/\sigma(e^\pm_R) \) for different weak interaction models. As you would expect, the SU(2) \(_L\) \( \otimes \) SU(2) \(_R\) \( \otimes \) U(1) model\( ^{55} \) gives considerably less parity violation than the SU(2) \(_L\) \( \otimes \) U(1) model\( ^{52} \). On the other hand, the right-handed doublet model\( ^{51} \) (I.61) gives more parity violation than the SU(2) \(_L\) \( \otimes \) U(1) model. Figure I.25 shows the effect on \( \sigma(e^\pm_L)/\sigma(e^\pm_R) \) of varying \( m_Z \). We have kept the same kinematic conditions as before, kept \( \sin^2 \theta_W = 0.25 \), and varied \( m_Z \) from the value (I.57) to \( \infty \), taking care to keep neutral current
cross-sections the same at $Q^2 = 0$. We see that parity violation effects would be larger if $m_z$ were much greater than the $O(80)$ GeV suggested by (1.57), and that it should therefore be possible to see $m_z$ propagator effects if they exist. Figure 1.26 compares $\sigma(e^+_l)/\sigma(e^-_R)$ at $x = 0.25$ for different values of $s$, using the Weinberg-Salam model with $\sin^2 \theta_W = 0.25$. We see that the CEEHP kinematic range lies above the threshold for seeing substantial parity violation effects, while there is not much to be gained by going to higher values of $s$.

All the results mentioned above were obtained in a simple-minded valence quark model where

$$u = 2d, \quad \bar{u} = \bar{d} = s = s = c = c = \ldots = 0$$

at $x = 0.25$, \hspace{1cm} (1.62)

and one may wonder how sensitive these results are to the possibility of scaling violations. Most models that we know of predict a similar dominance of valence quarks for $x > 0.2$, and we expect the previous results will not be greatly changed. As an example, the asymptotically free quark-parton distributions (1.14) and (1.15) have been used in Fig. 1.27 in a comparison with the valence model (1.62). We see that there is very little change in the calculated parity violation in the ratio $\sigma(e^+_l)/\sigma(e^-_R)$. Reference back to Fig. 1.22 shows that charge asymmetries $[\sigma(e^+_l, R) \neq \sigma(e^-_l, R)]$ can also be substantial. As an example, Fig. 1.28 shows the ratio $[\sigma(e^+_l) + \sigma(e^-_R)]/\sigma(e^+_l) + \sigma(e^-_R)]$ for the different weak interaction models discussed earlier. We see ratios substantially different from 1.

One may wonder whether neutral weak effects can be separated from multiphoton exchanges. The leading correction due to two-photon exchange comes from the interference diagram of Fig. 1.29a and has been estimated within the quark parton model. Assuming that both photons couple to the same parton line, we find that for electron scattering at large $Q^2$,

$$\frac{\sigma(2\gamma)}{\sigma(\gamma\gamma)} = -\frac{3}{2} \frac{a}{2\pi} Q_p^2 \ln^2 \frac{Q^2}{\mu^2},$$

(1.63)

where $Q_p$ is the charge of the quark parton ($\gamma_5$ or $-\gamma_5$), and $M$ is some scale mass, for which the most natural candidate is the parton mass. In order to determine the size of $M$ from experiment, we can compare Eq. (1.63) with the measurements of $\sigma(e^+ + p \rightarrow p + e^+ + X)/\sigma(e^- + p \rightarrow e^- + X)$ up to $Q^2 \sim 15$ GeV$^2$ and find that $M$ cannot be smaller than 0.3 GeV. At $Q^2 \sim 10^4$ GeV$^2$ the two-photon exchange contribution could then be as large as 20% of the one-photon exchange. When averaging over electron and positron scattering the interference diagram Fig. 1.29a drops out, and the leading corrections come from Fig. 1.29b. Their size has also been estimated, and for $Q^2 \sim 10^4$ GeV$^2$ does not exceed 2-3% of the one-photon exchange. In addition to these multiphoton exchanges, we must take into account radiative corrections at the leptonic vertex, which are much larger than other QED effects because of the smallness of the electron mass, and bremsstrahlung from both leptonic and hadronic parts. These processes all lead to terms $O(a/2\pi \ln^2 Q^2)$. The relatively slow $Q^2$ dependence of these higher-order QED effects should enable them to be distinguished from the $z^2$ exchange effects in which we are primarily interested. These exhibit a much sharper variation with $Q^2$, as seen in Fig. 1.28. A possible tactic may be to estimate the $2\gamma$ asymmetry from data.
at relatively low $Q^2$, where $Z^0$ effects should be small and extract the scale mass $M$ in (1.63). We may then use the relatively slow $Q^2$ dependence of $\sigma(2\gamma)/\sigma(1\gamma)$ to give an upper bound on $\sigma(2\gamma)$ at larger values of $Q^2$, where neutral current asymmetries should be much larger. The dashed lines in Fig. I.28 indicate the uncertainties in $\sigma(e^-)/\sigma(e^+)\gamma$ exchange contributions treated in this way.

Finally, Fig. I.30 shows that if one compares $\sigma(e^-) + \sigma(e^+)$ to the point-like single-photon exchange contribution, the models studied show surprisingly little deviation from unity over all the kinematic range. This fact may be useful in the separation of different sources of scaling violation. As an example of another source of scaling violation, we have shown in Fig. I.30 a calculation based on asymptotic freedom, which gives a very different pattern of scaling violation as a function of $Q^2$.

The main conclusions of this section are that substantial parity violation and charge asymmetries may be expected at CEEP, and that these may be used to distinguish different weak interaction models. To exploit these effects fully would need polarized beams of both electrons and positrons. However, going to energies much larger than those available at CEEP does not seem to greatly increase the observable effects.

### 3.2 Charged current effects

The kinematic range available to CEEP allows the reaction $e + p \to \nu + X$ to be used to probe the conventional charged weak interactions at $\sqrt{Q^2}$ comparable to their natural mass-scale of 50 to 100 GeV, corresponding to a neutrino beam of up to 14 TeV. Figure I.3 shows that there are usable event rates at $Q^2 = O(10^5)$, and the analysis given in Chapter II shows that these events are experimentally distinguishable. It will be of fundamental significance to see if there are deviations from the simple point-like four-fermion interaction, such as a damping from unitarity or $W$-boson propagator effects. The space-time structure of the charged current can also be probed to see if it is really pure $(V-A)$, and we can imagine the single production of new quark or lepton flavours with masses up to $O(50)$ GeV.

The weak cross-section can be written in terms of three structure functions $F_i(x, Q^2)$:

$$\frac{d^2\sigma}{dx dy} (e + p \to \nu + X) = \frac{G_F^2}{2\pi} s \bar{P}(s, x, y) \left[ (1-y)F_2(x, Q^2) + y^2 x F_1(x, Q^2) - y \left(1 - \frac{y}{x}\right) x F_3(x, Q^2) \right]$$

(1.64)

where we have allowed for the possibility of a deviation $\bar{P}(Q^2)$ from a point-like interaction: as examples,

$$\bar{P}(Q^2) = \begin{cases} 1 & \text{for a four-fermion theory} \\ \frac{m_W^2}{m_W^2 + Q^2} \times \text{for a single } W \text{ exchange} \end{cases}$$

(1.65)

The exact scaling curves in Fig. I.3 were calculated with the usual naïve scaling laws:

$$F_i(x, Q^2) = F_i(x)$$

(1.66)
and assuming the Callan-Gross relation\(^{19}\) shown in equation (1.3),

\[
F_2(x) = 2x F_1(x) ,
\]

and using the Barger and Phillips\(^{11}\) parametrization of the quark-parton distributions for the structure functions (1.66). The numbered curves\(^2\) in Fig. I.3 correspond to different possibilities for the W boson mass in Eq. (1.65). We see that at large \(Q^2\) a W boson mass of 62 GeV would reduce the cross-section by more than an order of magnitude, while even a W mass of 150 GeV should be easily distinguishable from a point-like weak interaction (\(m_W = \infty\)). This point is also brought out in Fig. I.31, where the total cross-sections for \(e^- p \rightarrow \nu + X\) and \(e^+ p \rightarrow \overline{\nu} + X\) divided by \(s\) are plotted. At CHEEP energies the cross-section for \(m_W = 150\) differs by about \(\frac{1}{4}\) from a linear extrapolation with \(s\) of the lower energy cross-section.

It might be objected that the effect of a damping of the four-fermion interaction could be mimicked or masked by scaling violations at the hadronic vertex. For example, a parton form factor of the type (1.14) would have an effect similar to the form of \(P(Q^2)\) in (1.65). However, any scaling violations would also occur in the electromagnetic/neural current reaction \(e^- + p \rightarrow e^- + X\), so that comparing this with \(e^- + p \rightarrow \nu + X\) should reveal unambiguously something like a W propagator, if one assumes that scaling is violated similarly in weak and electromagnetic interactions, as is strongly suggested by the CVC hypothesis. This assumption is compatible with present analyses of scaling violations in \(\pi N\), \(N\), and \(\pi N\) scattering\(^{12,15}\), and is expected in all known models of deep inelastic scattering. The strong interactions of quark-partons are not expected to depend on the method by which they are given a large momentum kick. Furthermore, specific models of scaling violations have very different characteristics from the damping (1.65). For example, in a field theory, as discussed in Section 2 of this chapter, the structure functions \(\pi W\) and \(W\) should rise as \(Q^2 \rightarrow \infty\) at small \(x\), and fall only for large \(x\). These different characteristics are reflected in the dashed curves\(^2\) in Figs. I.3 and I.31, which represent calculations using asymptotic freedom. In this model, the scaling violations are larger at small \(s\) than at high \(s\), while the W propagator effect is more pronounced at high \(s\).

On the basis of these arguments we feel that distinguishing charged weak cross-section damping from plausible forms of scaling violation should be possible. Problems could only arise if the scaling violations in weak and electromagnetic interactions were different, or were much wider than in the models discussed in Section 2 of this chapter.

Another aspect of the charged weak interaction which can be studied is its space-time structure. The weak interaction between the electron and its neutrino is believed to be pure (V-A), in which case the cross-sections for \(e^- + p \rightarrow \nu + X\) and \(e^+ + p \rightarrow \overline{\nu} + X\) should be identically zero. This can be tested using the variable polarization of the lepton beam. The next section describes how we might look for new massive leptons coupled to the electron via left- or right-handed currents.

3.3 New particle production

We should distinguish three different types of new particles which may be produced by CHEEP: new quarks\(^{19}\), heavy leptons, and intermediate vector bosons. The presence of any sort of new particle can only increase the charged current cross-section compared to that
estimated with the conventional weak interactions and shown in Figs. I.23 and I.31. We should perhaps note that if the naive interpretation of present atomic physics parity violation experiments is correct, and the standard Weinberg-Salam model is wrong, either a new neutral lepton $E^0$ or new right-handed vector bosons $W_R^+$, or both, very probably exist.

3.3.1 New quarks

Since CHEEP has a centre-of-mass energy $O(5)$ higher than present-day neutrino beams, it can look for new quark flavours with masses up to five times higher. The rates of new quark production are sensitive to the $V-A$ nature of their weak couplings as well as the intrinsic strengths of their weak interactions, but are potentially very dramatic. For example, in a valence quark approximation a new weak doublet

$$\begin{bmatrix} t \\ d \end{bmatrix}_R$$

would increase the cross-section for $e^+ + p \rightarrow \bar{\nu} + X$ by a factor of up to 4, while a doublet

$$\begin{bmatrix} Q = \frac{5}{2} \\ u \end{bmatrix}_R$$

would increase the cross-section by a factor of up to 7. These would correspond to event rates of 250 or 500 per day, even if $m_W = 62$ GeV and there are scaling violations due to asymptotic freedom. The precise increase in cross-section would depend on the mass of the new quark, but even if $m_L$ were 40 GeV the increase given by (I.68) would be at least a factor of 2. Hence CHEEP has the potential to search for new quarks in a mass range inaccessible to PETRA or PEP. New quarks such as those in (I.68) or (I.69) would be produced in the wide-angle parton jet, so that their effects would not be lost down the beam pipe. New massive quarks are expected to have substantial semileptonic decay branching ratios $O(10 \text{ to } 20)\%$. A quark of mass $O(50)$ GeV might be expected to yield charged leptons with transverse momenta $O(10)$ GeV.

3.3.2 New leptons

Many weak interaction models have heavy neutral leptons which couple to the electron through the known weak currents. Examples include the doublet

$$\begin{bmatrix} E^+ \\ \nu_e \cos \alpha + E^0 \sin \alpha \\ e^- \end{bmatrix}_L$$

and the triplet

$$\begin{bmatrix} e^+ \\ \nu_e \cos \alpha + E^0 \sin \alpha \\ e^- \end{bmatrix}_L$$

For $s \gg m_L^2$, these leptons are produced with rates
\[
\frac{\sigma(e_R^- + p \rightarrow E^\pm + X)}{\sigma(e_L^- + p \rightarrow e^- + X)} \approx \frac{1}{3} \left( \frac{g_{e R L}}{g_{e L L e}} \right)^2 = \frac{1}{3}
\] (I.72)

\[
\frac{\sigma(e_L^- + p \rightarrow e^- + X)}{\sigma(e_L^- + p \rightarrow e^- + X)} = \left( \frac{g_{e L L}}{g_{e L L e}} \right)^2 = \tan^2 \alpha
\] (I.73)

In a simple quark-parton model these ratios are correct to within a factor of 2 for \( \sqrt{s} > 4m_L \). Hence even if \( m_W = 62 \) and the cross-sections are further depressed by the scaling violations typical of asymptotic freedom, the model (I.70) would yield \( > 60 \) E\( ^5 \) per day at \( s = 27,000 \) GeV\(^2 \) for \( m_{E^5} \) up to 35 GeV. Even if a heavy lepton mass were 100 GeV, there could be several events per day.

The models (I.70) and (I.71) can be regarded as rather conservative. For example, doublets akin to (I.70) would probably also exist for the \( \mu^- \) and \( \tau^- \). The associated \( M^5 \) and \( T^6 \) neutral leptons would then in general mix with the \( E^5 \). Therefore they could also be produced by the electron beam, although with reduced rates. Also, many people suggest models with more than one \( W^5 \) boson. One example is based on the group \( SU(2)_L \otimes SU(2)_R \otimes U(1) \), and has doublets like (I.70) coupling mainly to a new \( W^{5'} \) boson with mass \( \geq m_W \). In this model the production of new leptons would be accompanied predominantly by the production of new quarks, because the \( W^{5'} \) bosons couple mainly to new right-handed quark doublets like

\[
\begin{pmatrix}
    u_L \\
    b_L \\
    d_R \\
    t_R
\end{pmatrix},
\begin{pmatrix}
    t_L \\
    d_R \\
    s_R
\end{pmatrix}, \ldots
\] (I.74)

Although couplings like (I.74) have not yet shown up in neutrino scattering\(^{61}\), they may well appear at CHeP if the new quarks have masses between 10 and 50 GeV, or if they are mediated by a \( W^{5'} \) which has only V+A couplings and therefore decouples from \( \nu_L(\tilde{v}_R) \). Another model with a proliferation of new quarks and leptons is based on the group \( SU(3) \otimes U(1) \). It has triplets

\[
\begin{pmatrix}
    d_L \\
    u_L \\
    b_L
\end{pmatrix}, \begin{pmatrix}
    s_L \\
    c_L \\
    t_L
\end{pmatrix}, \ldots,
\begin{pmatrix}
    e^- \\
    \nu_e \\
    \mu^-
\end{pmatrix}_L, \begin{pmatrix}
    \nu^\tau \\
    \tau^-
\end{pmatrix}_L,
\] (I.75)

\[
\begin{pmatrix}
    b'^L \\
    t'^L \\
    s_R
\end{pmatrix}, \begin{pmatrix}
    t' \\
    t'^R \\
    s_R
\end{pmatrix}, \ldots,
\begin{pmatrix}
    e^- \\
    \nu^e \\
    \mu^-
\end{pmatrix}_R, \begin{pmatrix}
    \nu^\tau \\
    \tau^-
\end{pmatrix}_R,
\] (I.75)

which contain new charged (as well as neutral) leptons which are coupled to electrons by an off-diagonal neutral current. Once again, almost every new lepton event also contains a new quark.
The kinematics of heavy lepton events are again very favourable, because the lepton tends to emerge at wide angles, with its decay products forming a "lepton jet" on the opposite side of the beam axis from the "parton jet", as indicated in Fig. I.32 [except that leptons with masses $O(100)$ GeV will decay more isotropically]. Possible decay modes of heavy leptons include\(^{60}\)

\[
\begin{align*}
E^0 & \to e^- + \text{hadrons} + \nu_e + \mu^+ + \mu^- \\
& \to e^- + \mu^+ + \nu_\mu + \bar{\nu}_e \\
& \to e^- + e^+ + \nu_e \\
& \to e^- + \text{hadrons} + E^- + \nu_e \\
& \to \nu_e + \text{hadrons} + E^- + \mu^+ + \nu_\mu \\
& \to \nu_e + e^+ + e^- + E^- + e^+ + \nu_e
\end{align*}
\]

(1.76)

for neutral leptons, and

\[
\begin{align*}
E^- & \to e^- + \text{hadrons} + \nu_e + e^- + \bar{\nu}_e \\
& \to e^- + e^+ + e^- + \nu_e + \mu^- + \bar{\nu}_\mu \\
& \to e^- + \mu^+ + \mu^- + E^0 + \text{hadrons} \\
& \to e^- + \nu_e + \bar{\nu}_e + E^0 + e^- + \bar{\nu}_e \\
& \to \nu_e + \text{hadrons} + E^0 + \mu^- + \bar{\nu}_\mu
\end{align*}
\]

(1.77)

for charged leptons. Many of the decay modes in (1.76) and (1.77) are very distinctive: $e^+\mu$ events, three charged leptons $e^-\mu^+\mu^-$, whose invariant mass is $m_{E^-}$, $e^-$ hadron combinations with a mass $m_{E^0}$, and so on. The three-jet final-state structure with transverse momentum approximately balanced would provide a distinctive signature, so that it seems feasible to identify leptons with production rates of 1 or 2 per day. Only an $e^+e^-$ machine with a centre-of-mass energy bigger than 100 GeV could compete with CHEEP as a source of heavy leptons coupled to the electron, if they exist.

3.3.3 Intermediate vector bosons

Since CHEEP is apparently able to probe weak interaction cross-sections at sufficiently high momentum transfers to see deviations from a point-like four-fermion theory and detect vector boson propagator effects, we might expect it to be a copious source of physical $W^\pm$ or $Z^0$ bosons. This is unfortunately not the case: we estimate\(^{1}\) the total cross-sections for $W^\pm$ to $Z^0$ production to be $O(10^{-17})$ cm$^2$ at $s \sim 27,000$ GeV$^2$. Sizable fractions of these cross-sections, especially in the case of the $Z^0$, are expected to be relatively "clean", with a final state containing just a lepton, the vector boson, and a nucleon or $N^*$. However, the radiative Bethe-Heitler dilepton background is not obviously negligible in such events.

We will now investigate these questions in some detail using as a guide the minimal Weinberg-Salam\(^{52}\) model with just one charged $W^\pm$ pair, and a neutral $Z^0$. The relevant Feynman diagrams can be divided into two classes -- those with $W^\pm$ or $Z^0$ emission from the lepton vertex shown in Fig. I.33, and those with emission from the hadron vertex shown in Fig. I.34. In both cases there is an exchanged photon which should be quasi-real in order
to get a substantial cross-section. The various diagrams in the two classes are interrelated by gauge invariance, meaning for example that in the hadronic diagrams "Drell-Yan" diagrams like Fig. I.34b cannot be considered in isolation from "bremsstrahlung" diagrams like Fig. I.34a. Detailed calculations of the resulting cross-sections and distributions from leptons produced in $W$ or $Z_0$ decay are still in progress, but the following preliminary estimates of the cross-sections can be given.

As far as the $W^\pm$ are concerned, we expect\textsuperscript{2} that from the leptonic vertex

$$\sigma(e^\pm + p \to W^\pm + \nu + X)_{\text{lept.}} \approx \sigma(\mu^\pm + p \to W^\pm + \nu + X)_{\text{lept.}},$$

(I.78)

where the latter cross-section has been calculated\textsuperscript{63} for a range of $W$ masses and beam energies. Figure I.35 shows the expected excitation curve as a function of $s/m_W^2$. If we take $m_W \approx 65$ GeV and $s = 27,000$ GeV$^2$ for CHeEP, then we find

$$\sigma(e^\pm + p \to W^\pm + \nu + p)_{\text{lept.}} = 3 \times 10^{-39} \text{ cm}^2$$

(I.79)

$$\sigma(e^\pm + p \to W^\pm + \nu + X \neq p)_{\text{lept.}} = 5 \times 10^{-39} \text{ cm}^2$$

which would be increased by a factor of 2 for polarized $e^-_L$ or $e^+_R$ beams. As for the hadronic vertex, if we use\textsuperscript{2} a simple valence quark model for the proton, which is surely inaccurate but should give the correct order of magnitude, we expect

$$\sigma(e + p \to e + W^\pm + X)_{\text{had.}} \approx \frac{2}{9} \sigma(\nu_\mu + p \to \mu + W + X)$$

(I.80)

$$\sigma(e + p \to e + W^\pm + X)_{\text{had.}} \approx \frac{1}{9} \sigma(\nu_\mu + p \to \mu + W + X)$$

Using previous calculations\textsuperscript{63} of the right-hand side of equation (I.80), we estimate that for $m_{W^\pm} \approx 65$ GeV, $s = 27,000$ GeV$^2$,

$$\sigma(e + p \to e + W^\pm + X)_{\text{had.}} \approx 5 \times 10^{-38} \text{ cm}^2$$

(I.81)

$$\sigma(e + p \to e + W^\pm + X)_{\text{had.}} \approx 2 \times 10^{-38} \text{ cm}^2$$

Combining the estimates (I.79) and (I.81) we see that for $m_{W^\pm} = 65$ GeV, $s = 27,000$ GeV$^2$ the total cross-section

$$\sigma(e + p \to \nu + W^\pm + X) \approx 10^{-37} \text{ cm}^2.$$

(I.82)

About 10% of the events would be relatively "clean" with only a small excitation at the hadronic vertex: these would include most of the events corresponding to lepton vertex emission of the $W^\pm$, and perhaps some of the hadron vertex events. The shape of the excitation curve as a function of energy would probably be similar to the shape of Fig. I.35.

As far as the $Z_0$ is concerned, we expect\textsuperscript{2,64} that from the leptonic vertex
\[
\sigma(e+p\rightarrow e^{\pm}Z\pm X)_{\text{lep.}} = \left( \frac{g_{V}^{2} + g_{A}^{2}}{2\sqrt{\alpha_{e}}} \right) \ln \left[ \frac{m_{Z}^{2}}{35m_{W}^{2}} \left( \frac{s}{m_{Z}^{2}} - 1 \right) \right] \sigma(\nu\bar{\nu}p\rightarrow \nu\bar{\nu}W\pm X)_{\text{lep.}}, \quad (I.83)
\]

where \( g_{V} \) and \( g_{A} \) are defined by the coupling

\[
2^{\alpha} \lambda_{V} (g_{V}^{2} + g_{A}^{2}) e. \quad (I.84)
\]

If we take \( m_{Z} = 75 \text{ GeV} \) and \( s = 27,000 \text{ GeV}^{2} \), and use the excitation curves corresponding to (I.83) plotted in Fig. I.36 we find

\[
\begin{align*}
\sigma(e+p\rightarrow e^{\pm}Z\pm p)_{\text{lep.}} &\approx 8 \times 10^{-38} \text{ cm}^{2} \\
\sigma(e+p\rightarrow e^{\pm}Z\pm X\#p)_{\text{lep.}} &\approx 6 \times 10^{-38} \text{ cm}^{2}
\end{align*}
\]

As for the hadronic vertex, the simple valence quark model for the proton used earlier (I.80) to estimate \( W^{\pm} \) production suggests\(^{63}\)

\[
\sigma(e+p\rightarrow e^{\pm}Z^{\pm}X)_{\text{had.}} \approx 0.15 \sigma(\nu\bar{\nu}p\rightarrow \nu\bar{\nu}W\pm X) \quad (I.86)
\]

which, when combined with numerical estimates\(^{63}\) of the right-hand side, yields for \( m_{Z} = 75 \text{ GeV} \), \( s = 27,000 \text{ GeV}^{2} \),

\[
\sigma(e+p\rightarrow e^{\pm}Z^{\pm}X)_{\text{had.}} \approx 2 \times 10^{-38} \text{ cm}^{2}. \quad (I.87)
\]

Scanning the estimates (I.85) and (I.87) we see that in total

\[
\sigma(e+p\rightarrow e^{\pm}Z^{\pm}X) = 10^{-37} \text{ cm}^{2} \quad (I.88)
\]

for \( m_{Z} = 75 \text{ GeV} \) and \( s = 27,000 \text{ GeV}^{2} \). In this case, we see from (I.85) that at least a half of the events are probably "clean" with little or no excitation of the hadronic vertex. The total excitation curve as a function of \((s/m_{Z}^{2})\) probably has a similar shape to those shown in Fig. I.36.

The existence of "clean" events with little or no excitation of the proton raises the prospect of detecting the \( W^{\pm} \) or \( Z^{\pm} \) inclusively, without having to pay the price of a leptonic branching ratio:

\[
\begin{align*}
\frac{\Gamma(W^{\pm}\rightarrow e^{\pm}\nu \text{ or } \mu^{\pm}\nu)}{\Gamma(W^{\pm}\rightarrow \text{all})} &\approx \frac{1}{4\pi} \\
\frac{\Gamma(Z^{\pm}\rightarrow e^{\pm}e^{-} \text{ or } \mu^{\pm}\mu^{-})}{\Gamma(Z^{\pm}\rightarrow \text{all})} &\approx \frac{3}{32\pi} \text{ for } \sin^{2} \theta_{W} \approx 0.25
\end{align*}
\]

(I.85)
where $N_D$ is the total number of either lepton or (coloured) quark fermion doublets, which we assume to be equal. Present experiments already suggest $N_D \geq 3$ so that the branching ratios in (1.89) are $0(5$ to $10)\%$. A potential background to the search for leptonic decays of $W^\pm$ or $Z^0$ is indicated in Fig. I.37. The two-photon differential cross-section for lepton-pair production is

$$\frac{d\sigma}{d\Omega_{\mu+\mu^-}} = \frac{4\alpha^4}{m^2_{\mu+\mu^-}} \left[ \ln \left( \frac{s}{m^2_{\mu+\mu^-}} \right) - 1 \right] \left[ \ln \left( \frac{s}{m^2_{\mu+\mu^-}} \right) - \frac{1}{4} \right]$$  \hspace{1cm} (I.90)

which gives\(^6\)\(^5\) cross-sections which are not totally negligible:

$$\sigma(70 \text{ GeV} < m_{\mu^+\mu^-} < 80 \text{ GeV}) \approx 4 \times 10^{-39} \text{ cm}^2.$$  \hspace{1cm} (I.91)

Calculations are in progress\(^6\)\(^6\) of the lepton spectra from vector boson production and decay which should enable us to see whether the boson signal can be distinguished from the $2\gamma$ background (I.91).

It is clear from the discussion of this subsection that an $e^-p$ machine is not the best for manufacturing $W^\pm$ or $Z^0$ bosons. Their production is more copious, if less "clean", in hadron-hadron scattering, whereas detailed studies of the decay modes and couplings of $Z^0$ bosons will best be realized with an $e^+e^-$ machine which can be tuned to $2E_{\text{beam}} = m_{Z^0}$ \(^6\)\(^7\).

3.3.4 Exotic

In this subsection we will briefly mention some inhabitants of the theoretical zoo which could conceivably be produced in high-energy $e^-p$ collisions. Not all their production mechanisms go through the weak interactions, but it seems natural to discuss them together.

**Higgs bosons** are required in all spontaneously broken gauge theories of the weak and electromagnetic interactions. The Higgs sector may be relatively simple, as in the standard Weinberg-Salam\(^5\)\(^2\) model where there is just one neutral scalar boson $H^0$ with specified couplings

$$\begin{align*}
\mathcal{H} = \frac{\mu^2}{2m_W} : & \quad g^2 = 4\sqrt{2} \frac{m^2_W}{G_F} \\
\mathcal{H}^0 VV = \frac{\mu^2}{m_W} & \quad \text{for } V = W^\pm \text{ or } Z^0
\end{align*}$$  \hspace{1cm} (I.92)

but almost completely arbitrary mass. On the other hand, there may be a large number of both neutral and charged Higgs mesons, as in most other weak gauge models. The predominant mechanism for production of the Weinberg-Salam Higgs $H^0$ is expected to be that shown in Fig. I.38. Use of the couplings suggests\(^6\)\(^8\) that for $s \gg m^2_{H^0}$,
\[ \frac{\sigma(e^- + p + \nu_e + H^0 + X)}{\sigma(e^- + p + \nu_e + X)} \leq \frac{G_F(x) s}{12\sqrt{2}\pi^2} \approx 10^{-8} \quad \text{for} \quad s = 27,000 \text{ GeV}, \quad (1.93) \]

where the inequality refers to neglect of the extra propagator suppression. Similar orders of magnitude hold for \( e^+ + \bar{\nu}_e + H^0 \) and \( e^+ + e^- + H \) via \( Z \)-exchange. This cross-section probably also gives a reasonable order of magnitude estimate for the production of neutral Higgs particles in more complicated models. Clearly the cross-section is unappetizingly small. Furthermore, we can discern no obvious distinguishing kinematical feature of Higgs production events. Also, the couplings (1.92) suggest that the dominant decay modes of Higgs bosons with masses \( \lesssim 2m_W \) will be to pairs of the heaviest kinematically accessible fermions. Given the existence of at least one heavy quark, the branching ratio into a readily identifiable final state such as \( \mu^+\mu^- \) will be impossibly small:

\[ \frac{\Gamma(H^0 \rightarrow \mu^+\mu^-)}{\Gamma(H^0 \rightarrow \text{all})} \lesssim \frac{m^2_{\mu}}{3m^2_Q}, \quad m_Q \sim 5 \text{ GeV} \]

\[ \sim \mathcal{O}(10^{-5}) \quad . \]

The prospects for charged Higgs \( H^\pm \) production, if such objects exist, are less grim because analogues of the electromagnetic diagrams of Fig. I.33 and Fig. I.37 exist in this case. However, the couplings and hence decay modes are more difficult to predict in the case of charged Higgs particles.

In view of the preceding remarks, we do not regard \( e^-p \) collisions as the optimal places to search for Higgs particles. As discussed elsewhere\(^{67} \), \( e^+e^- \) collisions are probably best for this purpose.

**Magnetic monopoles** are widely anticipated in gauge theories\(^{69} \). Their prediction is most often made in the context of unified gauge theories of weak and electromagnetic interactions, in which case

\[ m^2_{\text{monopole}} \approx \mathcal{O}\left(\frac{1}{\alpha}\right) m^2_W \quad (1.95) \]

so that the monopole mass is \( \mathcal{O}(10^6) \) GeV. This range is clearly far beyond the range of any presently conceivable machine. However, speculative arguments have been made for lighter monopoles, and one possibility\(^{70} \) is that in equation (1.95) \( m_W \) should be replaced by some hadronic vector meson mass, in which case a monopole mass \( \mathcal{O}(50) \) GeV is possible.

The experimental implications of monopoles are very striking\(^{71} \). Their coupling to the electromagnetic field is \( \mathcal{O}(1/\alpha) \), so that conventional perturbation theory breaks down. Below threshold, virtual monopole-antimonopole pairs may be produced which could annihilate to yield many photons. Above threshold, the pairs may interact, annihilate, escape or decay, with the emission of many photons. A monopole would interact electromagnetically very strongly in any detector, and would behave very dramatically in a magnetic field, aligning and accelerating itself parallel to the field.
The existence of magnetic monopoles in the mass range accessible to CHEEP is rank speculation, but they would certainly be noticeable if they were indeed produced. On the other hand, they might be produced earlier in high-energy hadron-hadron colliding ring experiments. It would be irresponsible not to draw the reader's attention to any of the bizarre high-energy cosmic-ray experiments.

Free fractionally charged quarks are anathema to many theoreticians, but ideas circulate as to their possible existence and experimental signatures. One school of thought suggests that free fractionally charged quarks would have small interaction cross-sections in ordinary matter\(^2\), in which case the experimental detection of their fractional charges \(\pm 1/6, \pm 2/3\) should be relatively straightforward. A more recent suggestion\(^3\) is that free quarks would have large interactions with ordinary nuclear matter, and a big appetite for absorbing nucleons. In this case, a free quark might emerge from the production reaction with its charge augmented by one or more units by dint of absorbing mesons or baryons produced in association. On passing through a particle detector it could in any case absorb more nucleons, increasing both its charge and especially its mass, producing an unfittable track.

Clearly the above two possibilities should be borne in mind during the design of detectors for high-energy e-p collisions. The mechanism of quark confinement or liberation is still totally obscure. Nevertheless, it is definitely suspected that the best way to liberate a quark may be to give it a large momentum transfer kick. In the absence of a high-energy e^+ e^- machine, e-p colliding rings will provide the cleanest samples of events with momentum transfers up to \(10^6\) GeV², and should therefore be the best place to look for free quarks, if they exist.
REFERENCES


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27) It should be pointed out that in some "colour liberation" models (Ref. 26) there are intermediate vector bosons which could cancel out some fraction of these increases.


38) E.G. Floratos, CERN preprint TH 2261 (1976) and contribution to this study.


44) F.E. Close, contribution to this study.


47) A. Ali, CERN preprint TH 2411 (1977) and contribution to this study.


55) J.C. Pati and A. Salam, Ref. 26;

56) V.G. Gorslikov et al., Soviet J. Nuclear Phys. 6 (1967) 361;


J. Smith, private communication.


65) P. Kessler and J. Parisi, contribution to this study. See also V.M. Rudnev et al.,

66) K.J.F. Gaemers, Work in progress and contribution to this study.
Quasi-diffractive production mechanisms of the form $\gamma + p \rightarrow Z + p$ have also been estimated in a covariant parton model. The resulting contributions to $\sigma(e + p \rightarrow e + Z + X)$ seem to be less than $10^{-5} \text{cm}^2$. J. Bartels and C.T. Sachrajda, contribution to this study.

67) L. Camilleri et al., CERN 76-18 (1976).

68) J. Ellis and M.K. Gaillard, contribution to this study. See also


70) W. Troost and P. Vinciarelli, CERN preprint TH 2195 (1976).

71) D.A. Ross, contribution to this study.


Fig. I.1 Deep inelastic scattering in the parton model

Fig. I.2 Event rates for $e + p \rightarrow e' + X$ in the one-photon exchange approximation
Fig. I.3 Event rates for $\bar{e}_L + p \rightarrow \nu + X$

$$e_L p \rightarrow X$$
Events per day with $Q^2 > Q_0^2$

$s = 40000 \text{ GeV}^2$, $L = 10^{32} \text{ cm}^{-2} \text{ sec}^{-1}$

- Exact scaling
- Asymptotic freedom

Fig. I.4 As $Q^2$ increases, a quark (a) may be resolved into (b) a quark and a bremsstrahlung gluon

Fig. I.5 As $Q^2$ increases, a "gluon" (a) may be resolved into (b) a quark-antiquark pair
Fig. 1.6 Field theory suggests the structure function will fall at large \( x \) and rise at small \( x \).

Fig. 1.7 The \( Q^2 \) dependence of \( W_2(x,Q^2) \) for \( x = 0.5 \), comparing asymptotic freedom and other field theories. The fixed point \( g^* \) for the Abelian vector gluon model has been chosen so that its predictions closely resemble asymptotic freedom for \( Q^2 \leq 30 \text{ GeV}^2 \). Uncertainties in mass corrections to scaling may compromise comparisons at \( Q^2 \leq 10 \text{ GeV}^2 \).
Fig. 1.8 The behaviour of $F_2(x, Q^2)$ in asymptotic freedom

Fig. 1.9 The expectation that the deep inelastic final state should have small $p_T$ in the photon-proton centre-of-mass frame
Fig. I.10 Conjectured regions of the longitudinal rapidity plot of the hadronic final states in (a) $e^-p$ collisions, (b) $p-p$ collisions, (c) $p + p \rightarrow \mu^+\mu^- + X$, and (d) $e^+e^-$ collisions.

Fig. I.11 Examples of how there may be non-local compensation of strangeness, charm, or other charge.
Fig. I.12 A model suggesting that multiplicities in hadron-hadron collisions should be twice as large as those in $e^+e^-$ collisions at high energies.

Fig. I.13 Possible non-scaling behaviour of the quark fragmentation function $D(z,Q^2)$.

Fig. I.14 Possible three-jet final state due to gluon bremsstrahlung.

Fig. I.15 Large $p_T$ cross-section from gluon bremsstrahlung.
Fig. I.16 The possible fragmentation of a heavy quark (a) into heavy and light hadrons, and the final states from (b) a semileptonic decay and (c) a non-leptonic decay of the heavy hadron $H$.

Fig. I.17 Examples of the transverse momentum spectra of leptons produced in the semileptonic decays of a meson made from a quark of mass 5 GeV.
Fig. I.18 A parton diagram may dominate Compton scattering at large energy and $t$.

Fig. I.19 Diffractive production of a vector meson.

Fig. I.20 Primakoff production of an even charge conjugation state.

Fig. I.21 Possible $Z_1^0$ exchange contributions to $e + p \rightarrow e' + X$.

$s = 27.000 \text{ GeV}^2$
$x = 0.25$

Fig. I.22 Cross-sections for $e + p \rightarrow e' + X$, normalized to one-photon exchange, for different electron and positron polarization states in the Weinberg-Salam model with $\sin^2 \theta_W = 0.25$. 
Fig. 1.23 The parity violating ratio $\sigma(e_l^-)/\sigma(e_R^+)$ for different values of $\sin^2 \theta_W$ in the Weinberg-Salam model.

Fig. 1.24 $\sigma(e_l^-)/\sigma(e_R^+)$ for different weak interaction models.
Fig. I.25 $\sigma(e^-)/\sigma(e^-)$ for different values of $m_z$, in the Weinberg-Salam model with $\sin^2 \theta_W = 0.25$

Fig. I.26 $\sigma(e^-)/\sigma(e^-)$ at $x = 0.25$ and $y = 0.5$, for different values of the beam energies, in the Weinberg-Salam model with $\sin^2 \theta_W = 0.25$

Fig. I.27 The effect of asymptotic freedom on $\sigma(e^-)/\sigma(e^-)$ in the Weinberg-Salam model with $\sin^2 \theta_W = 0.25$. 
Fig. 1.28 $\sigma(e^-)/\sigma(e^+)$ for different weak interaction models

Fig. 1.29 a) Two-photon, and b) multiphoton contributions to $\sigma(e + p \rightarrow e' + X)$
Fig. 1.30 Scaling violations and $\sigma(e^-) + \sigma(e^+)$ relative to a scaling cross-section in different weak interaction models.

Fig. 1.31 a) $\sigma(e^- p \rightarrow \nu_e X)$, and b) $\sigma(e^+ p \rightarrow \bar{\nu}_e X)$ as functions of the beam energies.
Fig. I.32 Event with a heavy lepton jet

Fig. I.33 Feynman diagrams for emission from the lepton vertex of a) charged $W^\pm$ bosons, b) neutral $Z^0$ bosons

Fig. I.34 Feynman diagrams for emission from the hadron vertex of charged $W^\pm$ or neutral $Z^0$ bosons classifiable as a) bremsstrahlung and b) Drell-Yan fusion
Fig. I.35 $\sigma(e+p \rightarrow \nu + W + X)_{\text{lept}}$ as a function of $s/m_W^2$

Fig. I.36 $\sigma(e+p \rightarrow e + Z + X)_{\text{lept}}$ as a function of $s/m_Z^2$
Fig. I.37 A $2\gamma$ mechanism for lepton pair production

Fig. I.38 A diagram for neutral Higgs boson production
CHAPTER II

PHYSICS WITH e-p COLLIDING BEAMS

Members of the ECFA Study Group *)

and

Participants in the Experimental Subgroup **) at the ECFA Study Week, Rutherford Lab., September 1977


1. **INTRODUCTION**

In the CHEEP option\(^1\) we propose to collide 25 GeV electrons with 270 GeV protons. This yields a c.m.s. energy of 164 GeV, equivalent to an incident electron energy of 14,400 GeV on a stationary proton target. The maximum value of \(Q^2\) is 27,000 GeV\(^2\). Protons at 400 GeV are available in the SPS at a reduced duty cycle, and the electron energy can be raised to 30 GeV by increasing the amount of RF power. This would enable CHEEP to reach a c.m.s. energy of 219 GeV, equivalent to a stationary target accelerator of 25.6 TeV. The maximum value of \(Q^2\) is 48,000 GeV\(^2\).

The CHEEP option makes it possible to investigate the strong, the electromagnetic, and the weak interaction in a kinematic region where new phenomena are expected to occur. This chapter is concerned with the experimental exploitation of CHEEP and considers the following topics.

The proposed layout of an interaction region is discussed in Section 2, and the background rates expected from synchrotron radiation and beam losses are given. A method for determining the luminosity is also presented. Section 3 deals with measurements of the inclusive lepton cross-section at large values of \(Q^2\). The kinematics, which are common to deep inelastic charged and neutral current events, are treated in the first part; the experimental problems, however, are very different for the two classes of events and are discussed separately in the second part. The design of a small calorimeter-type detector which is suitable for both types of experiments is presented. Section 4 deals with the additional information obtained by observing the final state in more detail, e.g. the production of hadrons with new flavours or of new leptons; in general a complex detector with particle identification is required. Some of the properties which such a detector must have are briefly discussed. Photoproduction is dealt with in Section 5.

2. **ENVIRONMENT**

2.1 **Layout of the interaction region**

The beam trajectories in the interaction region, and the layout of the magnetic elements, the vacuum chamber and the collimators, follow from the physics requirements -- namely, highest possible luminosity, electron or positron beams both longitudinally polarized, minimum background, and a sufficient amount of free space for experiments -- and from the very special machine environment, using protons of the SPS for the e-p collision.

Some parameters obtained in the machine study are listed in Table II.1.

The layout of the magnetic elements is shown in Chapter II. (Fig. III.50). The immediate vicinity of the free area for experiments is filled with bending magnets for separating the beams beyond the 5 mrad crossing angle, and quadrupoles to obtain the necessary low-beta values in the interaction region.

Because of the sizes of the proton and the electron beams, the separating magnets had to be common to both the electron (positron) and the proton beam for half of their total length. Therefore the horizontal septum magnet is split into four units of equal length, two of which act on both beams and the other two are double-gap magnets where the field in the gap on the proton side provides the necessary orbit corrections of the proton beam. The orbits in these magnets are given in Figs. III.10 and III.11, indicating also the horizontal apertures.
Table II.1
Parameter list

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Luminosity (peak)</td>
<td>$0.5 \times 10^{32}$/cm$^2$ sec</td>
</tr>
<tr>
<td>Electron energy</td>
<td>25 GeV</td>
</tr>
<tr>
<td>Proton energy range</td>
<td>150-400 GeV</td>
</tr>
<tr>
<td>Highest proton energy for &quot;infinite&quot; storage</td>
<td>270 GeV</td>
</tr>
<tr>
<td>Number of bunches</td>
<td>60</td>
</tr>
<tr>
<td>Proton bunch length</td>
<td>$\sim 30$ cm</td>
</tr>
<tr>
<td>Electron bunch length</td>
<td>$\sim 3-5$ cm</td>
</tr>
<tr>
<td>&quot;Natural&quot; polarization time at 25 GeV with &quot;polarization kinks&quot;</td>
<td>1.7 h</td>
</tr>
<tr>
<td>Critical energy of the synchrotron radiation created near the intersection</td>
<td>0.5 h</td>
</tr>
<tr>
<td>Free space for experiments</td>
<td>68.4 keV</td>
</tr>
<tr>
<td></td>
<td>$\pm 5$ m</td>
</tr>
</tbody>
</table>

In Fig. III.21 a field plot in one of the double magnets is shown.

The horizontal separating magnets are also meant to act as calorimeters to measure the forward proton fragments. The magnet yokes will consist of plates of 3 mm thickness interleaved with 3 mm thick scintillator. The minimum gap height is 130 mm. The proton beam passes very close (at a distance of $\sim 15$ mm only) to the septum coil in the double-gap magnet. Therefore that coil and also the other coils in these magnets should be made of aluminium, if possible.

The next machine element downstream from the proton beam is the yoke of the electron quadrupole, which is a variant of a Panofsky-type quadrupole. The yoke of this quadrupole can also be made in sheets in the same way as the yokes of the separating magnets. Figure II.1 shows the front face of that quadrupole with a yoke closing around the proton beam.

Further downstream are the proton quadrupoles, which have very stringent field requirements and large apertures ($\sim 200$ mm $\Phi$); they cannot be used as calorimeters. However, still further downstream there is place to intercept very small angle fragments of less than 3 mrad scattering angle.

Access to the area is provided through the access shaft to the long straight section LSS 5 of the SPS. The available lifting capacity and the available access to the shaft in the auxiliary building on top of the access shaft are not sufficient for the installation of the experimental apparatus. Most of the experimental apparatus could, however, be transported to the interaction area via LSS 6, which has a powerful lift. The available space in the straight section is sufficient for the installation of the calorimeter detector described in Section 3.2.

However, the full physics potential of CHEEP can only be realized using several different types of detectors. It therefore seems advantageous to make the halls so large that one experiment can collect data while a second is being mounted and tested. The hall described in connection with the "proton-antiproton" project in the SPS (Fig. II.2) is adequate for this purpose.
2.2 Synchrotron radiation

The special geometry of the interaction area, with fairly strong bending magnets near by, requires very special means for absorbing the synchrotron radiation far away from the detectors. Fortunately there is only one electron (positron) beam in an e-p machine. Here only the radiation generated by the bending magnets is treated, since the photons created in the quadrupoles are less numerous and are absorbed by the collimators foreseen for the radiation from the bending magnets. In Fig. II.3 the spectra of the synchrotron radiation originating from the common and double-gap magnets C1, C2, S1, S2 (field strength 0.16 T) and from the horizontal and vertical bending magnets H3, V1, etc., further upstream (field strength 0.28 T) are shown. For the names of the machine elements, see Fig. III.50. Tables II.2 and II.3 give an idea of the numbers involved.

<table>
<thead>
<tr>
<th>Table II.2</th>
</tr>
</thead>
<tbody>
<tr>
<td>Magnets H3, V1, etc.; critical energy: 115 keV; power line density: 6.1 kW/m.</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Photon energy range (keV)</th>
<th>Average (keV)</th>
<th>Photons per m keV sec</th>
<th>Photons per m sec</th>
<th>Photons per sec</th>
</tr>
</thead>
<tbody>
<tr>
<td>20–60</td>
<td>40</td>
<td>7 × 10^{15}</td>
<td>3 × 10^{17}</td>
<td>8 × 10^{17}</td>
</tr>
<tr>
<td>60–140</td>
<td>100</td>
<td>2 × 10^{15}</td>
<td>2 × 10^{17}</td>
<td>5 × 10^{17}</td>
</tr>
<tr>
<td>140–260</td>
<td>200</td>
<td>4 × 10^{14}</td>
<td>5 × 10^{16}</td>
<td>1 × 10^{17}</td>
</tr>
<tr>
<td>260–420</td>
<td>340</td>
<td>8 × 10^{13}</td>
<td>1 × 10^{16}</td>
<td>4 × 10^{16}</td>
</tr>
<tr>
<td>420–780</td>
<td>600</td>
<td>8 × 10^{12}</td>
<td>3 × 10^{15}</td>
<td>9 × 10^{15}</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Table II.3</th>
</tr>
</thead>
<tbody>
<tr>
<td>Magnets C1, C2, S1, and S2 critical energy: 68.4 keV; power line density: 2.1 kW/m.</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Photon energy range (keV)</th>
<th>Average (keV)</th>
<th>Photons per m keV sec</th>
<th>Photons per m sec</th>
<th>Photons per sec</th>
</tr>
</thead>
<tbody>
<tr>
<td>20–60</td>
<td>40</td>
<td>4 × 10^{15}</td>
<td>2 × 10^{17}</td>
<td>1.2 × 10^{18}</td>
</tr>
<tr>
<td>60–140</td>
<td>100</td>
<td>5 × 10^{14}</td>
<td>4 × 10^{16}</td>
<td>3 × 10^{17}</td>
</tr>
<tr>
<td>140–260</td>
<td>200</td>
<td>8 × 10^{13}</td>
<td>1 × 10^{16}</td>
<td>7 × 10^{16}</td>
</tr>
<tr>
<td>260–420</td>
<td>340</td>
<td>1 × 10^{13}</td>
<td>2 × 10^{15}</td>
<td>1 × 10^{16}</td>
</tr>
<tr>
<td>420–780</td>
<td>600</td>
<td>1 × 10^{12}</td>
<td>4 × 10^{14}</td>
<td>3 × 10^{15}</td>
</tr>
</tbody>
</table>

In Fig. II.4 the layout of machine elements upstream of magnet S2 is indicated together with a number of collimators (3 to 8). Of the total radiation, 100% from V1 and ~ 60% from H3 scatters the first time on collimators 8 and 7; ~ 10% from H3 on collimator 6, ~ 6% on collimator 5, 5% on collimator 4, 3% on collimator 3, 5% on collimator 2, 5% on collimator 1, and the remaining radiation crosses through the intersection in a similar way to the radiation
from the magnets C1, ..., S2, as described below. The radiation scattered from collimators 3 to 8 cannot reach the vacuum chamber in the interaction region without scattering at least once more because of the curved beam trajectory (Fig. II.5) leading into the interaction area. However, of the 3% of the total radiation of magnet H3 absorbed by collimator 3, a few percent are scattered into the septum magnet S2 which serves as a forward calorimeter. Scattered again in the vacuum chamber and the septum coil, the radiation finally reaching the scintillator gives rise to a pulse corresponding to an energy of the order of 10 GeV. This pulse occurs about 70 nsec before a real signal can arrive. The same happens in the septum magnet S1, and to a smaller extent in the magnets C1 and C2.

The synchrotron radiation originating in the magnets C1, C2, S1, and S2, and the remaining 6% of magnet H3, cannot be absorbed upstream of the intersection and has to pass through the intersection to be absorbed downstream. This absorption has to be done in such a way that all radiation scatters at least twice before reaching the vacuum chamber in the interaction region. In Fig. II.5 the horizontal layout of the vacuum chamber is explained, and Fig. II.6 gives the vertical dimensions of the special part through which the synchrotron radiation passes. This scheme requires a larger gap for C1' and C2'. After the first scattering off the absorber at -10 m away from the intersection, only 0.1% of the scattered radiation leaves the opening at -5.5 m without having scattered another time. Owing to the thin plates introduced into the opening, which present a very small target to the primary radiation, all the radiation entering the interaction area without having scattered twice is focused back towards the magnet C1. The same is done in the vertical view, where the opening towards the crossing point is smaller than the opening further away. Therefore only radiation scattered at least twice reaches the vacuum chamber in the sensitive area of the detector.

The scattered photons also arrive about 50 nsec after the real event. Of the number of photons mentioned above, the following numbers of photons leave the absorber area at -5.5 m after having scattered once (Table II.4).

<table>
<thead>
<tr>
<th>Photon energy range (keV)</th>
<th>Photons/sec leaving at -5.5 m</th>
</tr>
</thead>
<tbody>
<tr>
<td>20- 60</td>
<td>$2.4 \times 10^{12}$</td>
</tr>
<tr>
<td>60-140</td>
<td>$1.2 \times 10^{12}$</td>
</tr>
<tr>
<td>140-260</td>
<td>$4 \times 10^{11}$</td>
</tr>
<tr>
<td>260-420</td>
<td>$9 \times 10^{10}$</td>
</tr>
<tr>
<td>420-780</td>
<td>$4 \times 10^{10}$</td>
</tr>
</tbody>
</table>

Table II.4

Number of back-scattered photons

If all of these photons were aimed at the calorimeter and magnet C1 they would give rise to a pulse corresponding to several GeV total energy visible in the calorimeter, taking into account another scattering in the vacuum chamber and the relative acceptance of iron and scintillator. The pulse would be approximately 50 nsec after a real event in the calorimeter.
The number of photons which may reach the sensitive wire chambers in the centre of the interaction area can be obtained in the following way. In order to be able to scatter into the central vacuum chamber, photons have to scatter near the "exit" at roughly -6 m, which corresponds to a reduction by a factor of 500 of the original number of photons scattered once. After the second scattering, taking again the available solid angle to scatter into the detector region, the following maximum numbers of photons remain (Table II.5):

<table>
<thead>
<tr>
<th>Photon energy range (keV)</th>
<th>Photons/sec</th>
<th>Photons/bunch</th>
</tr>
</thead>
<tbody>
<tr>
<td>20-60</td>
<td>$10^8$</td>
<td>40</td>
</tr>
<tr>
<td>60-140</td>
<td>$10^8$</td>
<td>40</td>
</tr>
<tr>
<td>140-260</td>
<td>$4 \times 10^6$</td>
<td>2</td>
</tr>
<tr>
<td>260-420</td>
<td>$2 \times 10^7$</td>
<td>8</td>
</tr>
<tr>
<td>420-780</td>
<td>$10^7$</td>
<td>4</td>
</tr>
</tbody>
</table>

Thus approximately 100 photons reach the vacuum chamber surrounding the wire chambers. In a normal detector the conversion efficiency for photons is of the order of 10%. Therefore less than, say, 10 random digitizations will be visible in the wire chambers per event. The synchrotron radiation represents no problem for a calorimeter-type device. It is also very likely that drift and proportional chambers can be operated close to the beam line.

It should be mentioned briefly that with the present layout of magnetic elements around the crossing point the direct interaction of the synchrotron radiation with the proton beam is greatly reduced as compared to the e-p option\(^7\) of the LSR. A total rate of the order of 1 kHz of Compton-scattered photons and of $e^+e^-$ pairs will occur. These particles will travel at small angles with respect to the proton beam.

2.3 Beam-gas background

2.3.1 The proton beam

A variety of measurements\(^3\) was performed in LSS 5 of the SPS to study the rates of beam-gas and beam-wall interactions with the SPS operating for physics, and with bunched beams circulating at fixed energy for more than 20 min in machine development. In Fig. II.7 the rates measured in a scintillation counter telescope during three consecutive SPS cycles are shown. The machine operated with an intermediate flat top at 200 GeV where about 20\% to 30\% of the protons were slowly extracted. Then there was subsequent acceleration to 400 GeV, and fast ejection of the remaining protons. A large increase in rate is visible at low energies, corresponding to protons lost during injection and trapping which varies from pulse to pulse. Otherwise the rate increases with energy and is proportional to the number of protons at fixed energy. Both are compatible with predominant nuclear beam-gas scattering. The over-all loss rate seemed to have very little influence on the rates measured in LSS 5. In Fig. II.8 the rate measured, normalized to a fixed number of protons with a fixed energy beam versus time, is plotted, and correspondingly the decay rate of the proton beam. For a large variation of decay rate there is very little change in background rate.
By creating and displacing a pressure bump upstream of the counter set-up, the "source-length" of the beam-gas background was measured and found to be of the order of 50 m. The visible length was longer at higher energies than at lower energies because of the higher primary multiplicity and the more pronounced peaking at small angles of the background events. However, no measurement was done at very small angles, and the source length for those events may even be longer than 50 m.

The multiplicity of the background events was found to be of the order of \( \approx 15 \) charged particles/event.

A detector of 1 m outer radius with a hole of 15 cm radius centred around the beam pipe, with an average vacuum of \( \approx 5 \times 10^{-9} \) Torr, would see at present a total event rate of \( 5 \times 10^5 \) sec for \( 10^{12} \) circulating protons. Taking the CHEEP proposal of \( 2 \times 10^{13} \) protons in 60 bunches, we find \( 1 \times 10^7 \) events/sec or 4 events/proton bunch in the interaction area arriving with the proton bunch.

However, there are several ways of reducing this rate:

i) Possible improvements to the vacuum are at present being studied, and a gain of a factor 10 to 100 in the vicinity of the interaction region seems to be possible. This would reduce the background rates by a factor of \( \geq 10 \).

ii) Only a few percent of the particles in a background event penetrate one metre of iron. Adding shielding upstream of the detector should therefore reduce the rates by more than a factor of 10.

iii) Adding collimators has considerably reduced the background observed in ISR experiments. The same method can also be used at the SPS.

iv) The raw trigger rate can also be reduced considerably by demanding that the particles emerge from the nominal interaction volume using a fast logic.

These improvements will presumably reduce the basic trigger rate by several orders of magnitude down to a level where a more sophisticated analysis of the event can take place.

2.5.2 The electron beam

The background associated with the electron beam is presumably less serious than the one associated with the proton beam. This is supported by the experience gained at SPEAR and DORIS. Lost electrons will emerge along the direction of the electron beam with energies of the order of the primary electron energy and will not be confused\(^1\) with electrons from deep inelastic processes except in a small region of phase space. Electrons travelling on unstable orbits can also be removed by collimators mounted upstream of the interaction point in a region of large beam amplitudes.

2.4 Measurement of the luminosity

The luminosity must be known both in order to extract the cross-section from a measured rate and to optimize the performance of the accelerator. In the first case an attempt is made to minimize systematic errors; whereas to optimize the performance, a fast response for luminosities well below the nominal value is needed. Since these are conflicting requirements, we propose to use two separate luminosity systems, one based on elastic e-p scattering, the other on single bremsstrahlung.
The cross-section for elastic e-p scattering can be written as

\[
\frac{d\sigma}{dq^2} = \frac{4\pi\alpha^2}{q^6} \cos^2 \theta/2 \left[ A(q^2) + B(q^2) \frac{m_p^2}{s} \tan^2 \theta/2 \right].
\]

Since in the kinematical region of interest, \( \cos^2 \theta/2 \approx 1 \) and \( m_p^2/s \ll 1 \),

\[
\frac{d\sigma}{dq^2} \approx \frac{4\pi\alpha^2}{q^6} A(q^2) = \frac{4\pi\alpha^2}{q^6} \left[ G_E^2(q^2) + \frac{q^2}{4m_p^2} G_M^2(q^2) \right];
\]

\( G_E \) and \( G_M \) are the electric and the magnetic form factors, respectively. To estimate the rates, we use the approximate relations

\[
G_E(q^2) = G_M(q^2)/\mu,
\]

and

\[
G_B(q^2) = \frac{1}{(1 + 1.41 q^2)^2} \quad \text{with } q^2 \text{ in units of GeV}^2.
\]

This yields a cross-section:

\[
\frac{d\sigma}{dq^2} = 270 \text{ (nb/GeV}^2) \times \frac{1 + 2.2 q^2}{q^6(1 + 1.4 q^2)^2}.
\]

A luminosity of \( 10^{12} \text{ cm}^{-2} \text{ sec}^{-1} \) and a minimum scattering angle of 20 mrad result in 25 events/sec for 25 GeV electrons, increasing to 450 events/sec for an incident electron energy of 10 GeV. Thus the error in the determination of the luminosity is not dominated by statistical but rather by systematic uncertainties. The largest uncertainty is caused by the rapid variation of the elastic cross-section with angle. This error is given by

\[
\frac{d\sigma}{d\theta} \approx -4 \frac{d\sigma}{\theta} \quad \text{for } q^2 \ll 1 \text{ GeV}^2,
\]

\[
\frac{d\sigma}{d\theta} \approx -10 \frac{d\sigma}{\theta} \quad \text{for } q^2 \gg 1 \text{ GeV}^2.
\]

Note that the beam divergences at the interaction point are quite large -- for the electrons we have \( \gamma' = 10^{-6} \text{ rad}, x' = 2.3 \times 10^{-6} \text{ rad}. \) This error can be decreased by increasing the scattering angle accepting a corresponding reduction in rate. Using this method, the luminosity can presumably be determined with an uncertainty of a few percent.

The major background is expected to result from electron scattering on the residual gas. With a pressure of \( 10^{-10} \text{ Torr of mainly } H_2 \) \( ^{st} \text{ in the interaction region, the estimated background is small}. \) This can be seen by comparing the number of nuclei in the gas, about \( 10^9/m, \) with the number of protons per bunch, \( 3 \times 10^{11}. \) Also note that the beam-gas background is produced continuously along the beam direction, unlike real e-p scattering which originates in the interaction point.

The cross-section for electron-proton bremsstrahlung is given by

\[
d\sigma = 2\alpha^2 \int \frac{dk}{k} E k \left[ \frac{E^2 + E^2}{E_E E_E} - \frac{1}{3} \right] \times \left[ 2 \ln \frac{2E_E}{m_e E} - 1 \right];
\]
E₀ and E are the energies of the incident and scattered electron, respectively; k the photon energy; mₑ the electron mass; and r₀ = 2.81 × 10⁻¹³ cm. To estimate the rates, we use the fact that (dσ/dk)k is approximately constant to obtain

\[ \text{d}σ = 4 \times 10^{-26} \text{ d}k/\text{k cm}^2 . \]

The expected rate for a luminosity of 10¹² cm⁻² sec⁻¹ is 4 × 10⁶ photons with energies above 1/3 kₚ max per second, or 1.5 events/crossing. The rate is very large, and this reaction can therefore be used to monitor the luminosity over a wide range. The background -- mainly from beam-gas bremsstrahlung -- is not expected to be a serious problem. The problem caused by the synchrotron radiation cannot be evaluated before a detailed design of the interaction region has been made.

3. **Inclusive Lepton Measurements at Large Q²**

The interaction between an electron and a proton can to first order be represented by the Feynman graph shown in Fig. II.9:

![Feynman diagram](image)

Fig. II.9

The exchange current can be neutral (electromagnetic or weak), resulting in a charged lepton at the upper vertex; or it can be charged, leading to a neutrino or new neutral lepton. At the large values of Q² accessible with CHEEP, all these processes are of similar strength. Here we discuss the measurement of the lepton only, or equivalently, a measurement of the sum of all hadrons produced. We begin by reviewing the kinematics, which presumably are common to deep inelastic neutral and charged current events, and treat the experimental problems, which are different, separately.

3.1 **Kinematics**

The kinematical variables are defined above in Fig. II.9. The four-momenta of the incoming electron and proton and the outgoing lepton are denoted by e, p, and L, respectively; Q is the four-momentum of the current and W is the effective mass of the final hadronic system.

The invariants s, Q², and W² can be expressed in terms of these variables as:

\[ s = 4E_e E_p + m_e^2 \]
\[ Q^2 = -q^2 = 4E_e E_L \sin^2 \frac{\theta}{2} \]
\[ W^2 = Q^2 \left( \frac{E_p}{E_e} - 1 \right) + 4E_p(E_e - E_L) + m_p^2 \]
\[ = 2m_p \nu + m_p^2 - Q^2 \]
\[ \nu = \frac{p^a q}{m_p} = \frac{E_p}{2E_e} Q^2 + 2E_p(E_e - E_L) \frac{1}{m_p} \]
\[ = \frac{2E_p}{m_p} E_e - E_L \cos^2 \frac{\theta}{2}. \]

It is sometimes more advantageous to use the scale-independent variables
\[ x = \frac{Q^2}{2m_p \nu} \quad \text{and} \quad y = \frac{\nu}{\nu_{\max}} \quad \text{with} \quad \nu_{\max} = \frac{2E_p}{m_p}. \]

Table II.6 lists the formulas for the most important partial derivatives of \( Q^2 \), \( \nu \), and \( x \) with respect to the energy and angle of the outgoing lepton or the produced hadron system \( h \).

### Table II.6

<table>
<thead>
<tr>
<th>Partial Derivative</th>
<th>Formula</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \frac{\partial \nu}{\partial E_L} )</td>
<td>( \frac{\nu}{E_L} - \frac{2E_pE^p}{E_Lm_p} )</td>
</tr>
<tr>
<td>( \frac{\partial \nu}{\partial \theta_h} )</td>
<td>( \frac{Q^2}{\tan \theta/2} )</td>
</tr>
<tr>
<td>( \frac{\partial x}{\partial \theta_h} )</td>
<td>( \frac{x}{\tan \theta/2} \left( 1 - x \frac{E_p}{E_e} \right) )</td>
</tr>
<tr>
<td>( \frac{\partial \nu}{\partial \theta_h} )</td>
<td>( \frac{</td>
</tr>
<tr>
<td>( \frac{\partial \nu}{\partial p^\nu_h} )</td>
<td>( \frac{E_p}{m_p} )</td>
</tr>
<tr>
<td>( \frac{\partial x}{\partial \theta_h} )</td>
<td>( \frac{</td>
</tr>
<tr>
<td>( \frac{\partial x}{\partial \theta_h} )</td>
<td>( \frac{</td>
</tr>
</tbody>
</table>
\begin{table}[h]
\centering
\caption{Kinematics of the scattered electron and the current jet (all energies in GeV, $Q^2$ in GeV$^2$ and angles in degrees)}
\begin{tabular}{|c|c|c|c|c|c|}
\hline
\textbf{x} & \textbf{$y$ (GeV)} & \textbf{0.1 (1466)} & \textbf{0.3 (4379)} & \textbf{0.5 (7328)} & \textbf{0.7 (10,000)} & \textbf{0.9 (13,190)} \\
\hline
0.01 & $Q^2$ & 27.5 & 82.5 & 137 & 192 & 247 \\
 & $E'_e$ & 22.8 & 18.4 & 13.8 & 94 & 5 \\
 & $\theta_e$ & 12.6 & 24.5 & 37 & 53 & 85 \\
 & $\theta_J$ & 95.6 $\pm$ 10 & 54 $\pm$ 10 & 36 $\pm$ 7 & 23 $\pm$ 6 & 12 $\pm$ 3 \\
0.05 & $Q^2$ & 137 & 41.2 & 68.7 & 96.2 & 1238 \\
 & $E'_e$ & 23.9 & 21.6 & 19.4 & 17 & 14.9 \\
 & $\theta_e$ & 28 & 51.5 & 73 & 97 & 131 \\
 & $\theta_J$ & 131 $\pm$ 5 & 97 $\pm$ 5 & 72 $\pm$ 5 & 52 $\pm$ 5 & 26 $\pm$ 5 \\
0.1 & $Q^2$ & 275 & 82.5 & 1925 & 2475 \\
 & $E'_e$ & 25.2 & 25.7 & 26 & 27 \\
 & $\theta_e$ & 38 & 69 & 115 & 143 \\
 & $\theta_J$ & 144 $\pm$ 4 & 115 $\pm$ 4 & 68 $\pm$ 4 & 38 $\pm$ 2 \\
0.3 & $Q^2$ & 825 & 2475 & 1425 & 5775 & 7245 \\
 & $E'_e$ & 30.7 & 42 & 54 & 65 & 77 \\
 & $\theta_e$ & 62.4 & 97 & 122 & 140 & 159 \\
 & $\theta_J$ & 158 $\pm$ 4 & 138 $\pm$ 3 & 121 $\pm$ 3 & 99 $\pm$ 3 & 63 $\pm$ 3 \\
0.5 & $Q^2$ & 1375 & 4125 & 6875 & 9625 & 12,380 \\
 & $E'_e$ & 36 & 59 & 81 & 104 & 126 \\
 & $\theta_e$ & 76.2 & 114 & 133 & 149 & 163 \\
 & $\theta_J$ & 163 $\pm$ 2 & 148 $\pm$ 2 & 133 $\pm$ 3 & 113 $\pm$ 3 & 75 $\pm$ 3 \\
0.7 & $Q^2$ & 1925 & 5775 & 9625 & 13,480 & 17,330 \\
 & $E'_e$ & 41.8 & 75 & 108 & 142 & 175 \\
 & $\theta_e$ & 85 & 122 & 140 & 152 & 166 \\
 & $\theta_J$ & 166 $\pm$ 2 & 153 $\pm$ 2 & 139 $\pm$ 2 & 122 $\pm$ 3 & 83 $\pm$ 3 \\
0.9 & $Q^2$ & 2475 & 7425 & 12,486 & 17,330 & 22,280 \\
 & $E'_e$ & 47 & 91 & 136 & 181 & 225 \\
 & $\theta_e$ & 92 & 127 & 144 & 156 & 167 \\
 & $\theta_J$ & 167 $\pm$ 2 & 156 $\pm$ 2 & 144 $\pm$ 2 & 128 $\pm$ 2 & 92 $\pm$ 2 \\
\hline
\end{tabular}
\end{table}

Figure II.10 shows the kinematical boundaries in the $\nu$, $Q^2$ plane for 25 GeV electrons on 270 GeV protons. The solid line represents $W = \text{constant}$, the dashed lines show constant scattering angle ($\theta_L$) and lepton energy ($E_L$). Note that the lepton generally appears at large scattering angles with high energy and is therefore easy to identify and measure.

The kinematic of the scattered electron is shown in a polar diagram ($x = p^L_T, y = p^L_T$) in Fig. II.11. The incident electron and proton move along the positive and negative $x$-axis, respectively; the interaction point is at the origin. The curves of constant $W$ (ellipses)...
and of constant $Q^2$ (parabolas) are shown. The lepton is scattered into the forward hemisphere for small values of $Q^2$ only; the angle between the scattered electron and its initial direction is always greater than $20^\circ$ for $Q^2 > 50$ GeV$^2$.

Only energy-momentum conservation was used to derive the electron kinematic, and the same would of course be true if we only wanted to know the kinematics of the sum of the hadrons. However, to obtain the hadron distribution in the final state in more detail, we must appeal to a model. The parton model, which is the only model known which is consistent with all experimental data available on deep inelastic processes, predicts a final state consisting of two jets of hadrons as shown schematically in Fig. II.12. In this model the current interacts with a light parton, which carries a well-defined fraction $x$ of the proton momentum and has a small transverse momentum.

After the collision the parton that is struck materializes into a jet of hadrons at large angles (current jet), while the proton jet (the remains of the proton) gives rise to a jet along the initial proton direction. Both jets are well separated from the scattered electron. The momentum of the current fragmentation jet ($x \bar{p}_T + \bar{q}$) is determined by the electron variables, i.e. the transverse momentum of the scattered lepton with respect to the beam axis is balanced by the transverse momentum of the current jet with the proton jet carrying zero net transverse momentum. The polar diagram of the current jet with lines of constant $W$ and $Q^2$ is shown in Fig. II.13. In Fig. II.14 the information is condensed into a single polar diagram of the outgoing electron, with the directions of the current and fragmentation jets shown. The scale of the momentum vectors has been reduced by a factor of 10 in order to make the scale readable.

To ascertain that the current jet and the proton jet are indeed well defined in space, events were produced assuming an inclusive hadron distribution of the form

$$f(x_F, p^2) = e^{-4p_T^2} e^{-0.75(n) x_F};$$

$(n)$ denotes the mean number of charged particles produced and is given by $(n) = 0.36 + 1.05 \ln k^2 - 0.15 \ln Q^2$. This ansatz is in agreement with present data and ensures that energy and momentum are conserved.

The results of Monte Carlo computations for $(x, y) = (0.3, 0.9), (0.9, 0.3)$, and $(0.5, 0.5)$ are shown in Fig. II.15 and Fig. II.16. The number of hadrons is plotted as a function of the production angle $\theta$ for various values of the parameters. In all cases the hadrons appear in two well-defined jets, one along the initial proton direction, the other (the current jet) at an angle determined by $x$ and $y$. The hadrons in the current jet are contained within an angular cone with an opening angle of roughly three degrees. This distribution widens at small values of $x$. This is quantitatively shown in Table II.7, where $\theta_j$ is given with its r.m.s. value obtained by the Monte Carlo calculations. It is important to note that this two-jet structure is present even for a completely flat rapidity distribution assuming only a limiting transverse distribution of the hadrons.

If the proton is at rest then the current jet and the scattered lepton will appear in the forward direction on opposite sides of the incident beam line. The "proton jet" (target fragmentation) produces slow particles distributed isotropically. Figure II.17 shows an interesting example of a deep inelastic muon event$^6$ displaying the expected structure. Note that the target fragments are not observed with this set-up.
3.2 The layout of a compact CHEEP detector

It is the aim of this section to describe a compact non-magnetic detector optimized for the measurement of inclusive charged and neutral electron-proton interactions. It will be seen that the detector, as designed, will fit in the existing SPS tunnel, and there would be no appreciable difference if there were unlimited space available. Also, although the detector is simple, it will probably make a better job of the inclusive measurements of both charged and neutral current events than would more elaborate detectors with magnetic momentum measurement and particle identification, which are vital for measurement of specific channels.

The detector is based on energy measurement by the total absorption or calorimeter technique. The design is aided considerably by the fact that the kinematics of the reactions determine the energies and angles to be expected for electrons and for hadrons, with the barest minimum of assumptions about the details of the hadron dynamics as discussed in Section 3.1. The situation is thus different from the one encountered in p-p collisions, for example. The kinematics determine the division of the calorimeter into four pieces, as shown in Fig. II.18. These are given names, for convenience:

i) HIREM, a high resolution calorimeter for electromagnetic showers covering an angular region where many electrons but few hadrons are expected;

ii) CENCAL, a calorimeter covering most of the solid angle, capable of measuring both electromagnetic and hadronic showers with moderately good spatial and angular resolution;

iii) FORC, a calorimeter designed to measure the energetic particles emerging in the forward direction with respect to the proton beam, with good energy resolution and the best feasible angular resolution;

iv) FRAGCALS, a combined bending magnet and calorimeter which measures the most forward particles, essentially fragments of the target proton, with moderate energy resolution and poor angular resolution.

The characteristics and functions of each of these detectors will be described successively and are summarized in Table II.8.

HIREM

This is a calorimeter covering angles up to 180 mrad with respect to the electron direction. This corresponds to $x > 0.9, y > 0.99$ in the hadron system, and we may safely assume that only a completely negligible amount of hadronic energy reaches this portion of phase space. We take advantage of this circumstance to provide a detector which is designed to give the best resolution for electrons. The electrons will lie in the range from 1 GeV to 25 GeV, concentrated at the upper end, corresponding to inelastic scattering with $Q^2 < 30 \text{ GeV}^2$, $W < 150 \text{ GeV}$. The electron angle is measured to adequate precision in multiwire proportional chambers (MWPCs) with a lever arm of 3 m, so that the error in $Q^2$ is determined by the energy measurement, as is the error on $W$. The detector proposed is a liquid argon (LA) ion chamber with 1500 plates of 0.1 mm copper separated by 1.5 mm LA gaps. The resolution of a calorimeter with thicker lead plates was measured to be $4\% \sqrt{E} \text{ r.m.s.}$ and the resolution of the configuration proposed here should be appreciably better. Allowing for systematic effects, the resolution should be $\leq 1\%$ over the range of interest. The first five radiation lengths are read out separately to allow a check of the electromagnetic character of the electron.
shower. The transverse subdivision of the detector can be as fine as a few millimetres without affecting the resolution, but a relatively coarse subdivision, say 100 cells, would seem to be adequate unless the analysis can be shown to benefit by the separate detection of the internally radiated photons.

A NaI scintillation calorimeter can also be considered for this detector with the aim of reaching a resolution of 0.25%. However, it might be difficult to maintain its calibration to that accuracy in the CHEEP environment, and probably it would be necessary to give up the subdivision.

CENCAL

This detector covers the central 160° and must measure scattered electrons as well as the wide-angle hadrons over a wide range of energy. It is essential to measure the hadrons well enough to give an accurate energy and momentum balance in order to identify the charged current events with unobserved neutrinos. We propose taking advantage of the better resolution obtained in a uranium calorimeter, which also gives a compact design owing to the short absorption length. About 33 cm of uranium are needed to absorb particles of 10 GeV. Read-out by liquid argon is suggested because it leads to a compact design with 1.25 mm gaps between the plates, which are 1.5 mm thick to avoid sampling fluctuation limitations. Thus the active depth of the calorimeter is less than 60 cm. The 90° portion of the calorimeter is divided into blocks, 10 cm square, with the first three and second three radiation lengths read out separately for electron identification. The portion of the calorimeter forward of 40° have the plates oriented normal to the beam direction with an equivalent arrangement of subdivisions. The electron energy resolution in this device will be considerably worse than in HIREM, about 10%/E^1 for electron and 50%/E^1 for hadrons. The longitudinal shower distribution in three regions should give a hadron rejection of several hundred, which should be quite sufficient to identify charged current events of the expected sort. More complex multi-lepton events may require a more elaborate detector.

With a calorimeter inner radius of 0.5 m, the detector will fit in the standard SPS tunnel, as shown in Fig. II.19. Even given more space, similar dimensions would probably be retained in order to keep the calorimeter mass at a minimum, since the hadron angular resolution seems to be adequate. The design foresees the division of the calorimeter into eight 15-ton modules which can be transported in the tunnel.

The calorimeter would be designed with an emphasis on reliability. For example, by reading out each section with independent, interleaved sets of plates, a high degree of redundancy is obtained for little additional cost.

The performance requirements on the MWPCs inside CENCAL are not very severe as far as accuracy is concerned, but it is desirable to maintain a good track pair resolution. The most important requirements are reliability and low mass in the parts through which the electrons must pass before reaching the calorimeter, particularly the end calorimeter. There is still not much experience with cylindrical chambers, and in a year or two it will be possible to make better decisions based on experience gained with the detectors at PETRA and PEP. The parameters in Table II.8 are typical of a possible design.
<table>
<thead>
<tr>
<th>HIRES</th>
<th>CENCAL</th>
<th>FORC</th>
<th>FRAGCAL</th>
</tr>
</thead>
<tbody>
<tr>
<td>Electromagnetic shower calorimeter</td>
<td>Hadronic and electromagnetic calorimeter</td>
<td>Hadronic and electromagnetic calorimeter</td>
<td>Hadronic calorimeter</td>
</tr>
<tr>
<td>Structure: 1.5 mm LA/0.1 mm Cu</td>
<td>Structure: 1.25 mm LA/1.5 mm $^{238}$U (Pb in first three r.l.)</td>
<td>Structure: 1.5 mm LA/3 mm Fe (Pb in first five r.l.)</td>
<td>Structure: 3 mm Fe/3 mm scintillator</td>
</tr>
<tr>
<td>Longitudinal division: 5 r.l./20 r.l.</td>
<td>Longitudinal division: 3 r.l./3 r.l./100 r.l.</td>
<td>Longitudinal divisions: 5 r.l./50 r.l.</td>
<td>Longitudinal division: none</td>
</tr>
<tr>
<td>Transverse divisions: $\phi = 30^\circ$, $\Delta \theta/\theta = 10%$</td>
<td>Transverse divisions: 10 cm $\times$ 10 cm</td>
<td>Transverse divisions: 5 cm $\times$ 5 cm</td>
<td>Transverse divisions: quadrants</td>
</tr>
<tr>
<td>Depth: 1.7 m</td>
<td>Depth: 60 cm</td>
<td>Depth: 1.5 m</td>
<td></td>
</tr>
<tr>
<td>Weight: 5 tons</td>
<td>Weight: 8 $\times$ 15 tons</td>
<td>Weight: 8 tons</td>
<td></td>
</tr>
<tr>
<td>Energy resolution: $3%/E_1^{1/2} + 0.2%$</td>
<td>Electron energy resolution: $10%/E_1^{1/2} + 0.5%$</td>
<td>Electron energy resolution: $10%/E_1^{1/2} + 0.5%$</td>
<td>Energy resolution: $60%/E_1^{1/2} + 2%$</td>
</tr>
</tbody>
</table>

Inner track detector, central cylindrical position:
- $r_{\text{min}} = 200$ mm
- $r_{\text{max}} = 410$ mm
- $L = 2$ m

Wire spacing: 2 mm

Number of radial layers: six with alternate right- and left-handed helical pitch, each with charge division to aid pattern recognition
FORC

A large fraction of the hadronic energy escapes through the aperture of CENCAL. Most of this is detected in the forward hadron calorimeter. Since the typical hadron energy is quite high, and there is not a stringent limitation on the length, iron plates should be adequate for this calorimeter giving a resolution of about $50\%/E^2$. The first five radiation lengths are read out separately to allow the identification of photons and the few primary electrons which reach this detector.

For this calorimeter and for CENCAL, a scintillator read-out could be used. The better time resolution which would result does not lead to an advantage in this application because of the time structure of the beams (and background). There would be operational advantages and a gain in reliability because of absence of inaccessible electrical connections. The disadvantages are the difficulty of maintaining the relative calibration of a large number of channels without the possibility of a beam line calibration, and the considerably lower density resulting from the larger inter-plate gaps necessary for the scintillators.

FRAGCAL

The kinematic considerations in Section 3.1 suggest that the remains of the proton will give rise to a jet of particles travelling at small angles with respect to the proton-beam direction. It is of crucial importance for the identification of charged current events that these particles be observed down to very small angles. The presence of the long, weak-field, electron bending magnet downstream from FORC provides an opportunity for catching these particles. The low field in this magnet allows it to be built of 3 mm iron plates interleaved with 3 mm scintillators. The scintillator is easily coupled to phototubes at the end of the magnet. The resolution should be roughly $60\%/E^2$. Each magnet is divided into quadrants for read-out. No longitudinal subdivision is provided.

3.3 Neutral currents

Inclusive electroproduction $e + p \rightarrow e' + X$ proceeds by both the electromagnetic and the neutral weak currents. The electromagnetic interaction dominates at small values of $Q^2$, but if the weak interaction remains point-like it would dominate for $Q^2$ larger than 5000 GeV$^2$. However, in present theories the neutral weak interaction is damped by assigning a mass of the order of 70 GeV to the conjectured $Z^0$. In this case one expects the electromagnetic and the weak interactions to be of similar strength for $Q^2 \geq m_{Z^0}^2$.

We will first evaluate the contribution from one-photon exchange. The contributions from higher order photon exchange are effectively cancelled by summing the cross-sections observed with incident electrons and positrons. This cross-section can be written as:

$$ \frac{d\sigma}{dx dy} = \frac{4\pi\alpha^2}{sx^2y^2} \left[ (1-y) F_2(x, Q^2) + y^2 \times F_1(x, Q^2) \right]. $$

We assume, in agreement with present data, that the transverse virtual photons dominate the longitudinal photons. This is expected in spin $\frac{1}{2}$ parton models and leads to

$$ F_2(x, Q^2) = 2 \times F_1(x, Q^2). $$
We further assume scaling, i.e. \( F_2(x, Q^2) = F_2(x) \). The cross-section in this case can be evaluated\(^{10,11}\) in terms of existing data. The resulting rates per day in bins of \( \Delta x = \Delta y = 0.04 \) are shown in Fig. II.20 for 25 GeV electrons on 270 GeV protons and a luminosity of \( 10^{32} \text{ cm}^{-2} \text{ sec}^{-1} \). In Fig. II.21 the rate of events with \( Q^2 > Q_0^2 \) are plotted\(^{10}\) as a function of \( Q^2 \).

These rates are expected to be modified by scaling violations, and the rate expected\(^{10,12}\) in an asymptotically free theory is also listed in Fig. II.20. The number in brackets is the rate expected in an asymptotically free theory divided by the rate obtained assuming scaling. This scaling violation leads to an increase in rate at low values of \( Q^2 \) and a decrease at large values of \( Q^2 \).

The electroproduction rate is large enough to exploit much of the available kinematical region provided the electrons are identified and measured over a large solid angle. In fact, because of the point-like nature of the interaction and the asymmetry between the electron and the proton energies, the scattered electron will in general emerge in the laboratory with high energy at rather large angles with respect to the direction of incidence. This can be seen by comparing Figs. II.10 and II.20. The non-magnetic detector discussed in Section 3.2 covers \( 2\pi \) in azimuthal angle and between \( 1^\circ \) and \( 179.5^\circ \) in production angles, i.e. data over the complete kinematical region can be collected at one setting.

We next consider the background resulting from the interaction of the beam particles with the residual gas. Although the beam-gas rate can be quite high, as discussed in Section 2.3, such events can only populate a small region of phase space. Proton-gas events, because of the strong Lorentz boost, will travel near to the incident proton direction and populate the narrow black area at large values of \( y \), as shown in Fig. II.22. Because of the low incident electron energy, interactions with the rest-gas can only populate the region of small \( Q^2 \) along the \( \nu \) axis, as shown in Fig. II.22. We therefore conclude that most of the kinematical region does not receive contributions from beam-gas interactions. Beam-gas interactions can be rejected off line, and only pose a problem as far as the trigger rate is concerned.

To identify the electron among the hadrons in a genuine deep inelastic event should be rather straightforward since the electron is quite energetic (Fig. II.10) and appears at large angles on the opposite side of the beam line with respect to the current jet, whereas the proton jet travels down the beam line.

Photoproduced pions, where a pion is wrongly identified as an electron, might cause problems in certain regions of phase space. The angular and the momentum distribution of inclusive charged pions in \( e^+ p \to e^- \pi^\pm + X \) has been estimated\(^3\) assuming Feynman scaling in \( x_F \).

The relevant expressions in the hadron c.m.s. are:

\[
E^* \frac{d\sigma}{dp^4} (\nu e + p \to \pi^\mp + X) = F(x_F, p_T^2),
\]

with

\[
F(x_F, p_T^2) \equiv F(x_F) \left\{ a \exp (-ap_T^2) + 0.1 \exp (-bp_T^2) \right\}
\]
and

\[ F(x_F) = \begin{align*}
20 \mu b \exp (6.3 x_F) & \quad \text{for } -1.00 \leq x_F < -0.05 \\
20 \mu b \exp (-x_F) & \quad \text{for } -0.05 \leq x_F < 0.05 \\
330 \mu b \exp (-5.4 x_F) & \quad \text{for } 0.05 \leq x_F < 1.00
\end{align*} \]

a = 8 \text{ GeV}^{-2}, \quad b = 2 \text{ GeV}^{-2}.

These computations\(^{13}\) show that the photoproduced pion background is only important in a very small region of phase space \((Q^2 \leq 100 \text{ GeV}^2, \nu \text{ large})\).

We therefore conclude that deep inelastic neutral current events have a unique signature with essentially no background from beam-gas interactions or photoproduction.

The resolution in \(Q^2\) and \(\nu\) for the proposed detector has been evaluated assuming that the energy and the angle of the electron are measured with an error \(dE_e/E_e = 0.1/\sqrt{E_e} \text{ (GeV)}\) and \(d\theta_e = 5 \text{ mrad}\). The result is shown in Fig. II.25. Plotted is \(d\nu/\nu\) and \(dQ^2/Q^2\) as a function of \(y\) for \(x = 0.1\) and 0.9. The resolution is clearly adequate for \(y > 0.2\).

Inclusive electron scattering at large values of \(Q^2\) provides a way of obtaining new information\(^{13}\) on the coupling of the neutral weak current to leptons and hadrons. It also allows a direct comparison between the electromagnetic and the weak interaction in a single reaction.

Experimentally there are several ways of identifying the weak contribution to \(e + p \rightarrow e' + X\).

i) **Parity violation.** The weak cross-section depends on the helicity of the incident lepton. This is a unique signature of the weak current.

ii) **The weak interaction is different for \(e^-\) and \(e^+\).** This can in principle be masked by two-photon effects. However, calculations show that the weak contribution is much larger than two-photon effects and has a different \(Q^2\) dependence.

iii) **Deviations from the Rosenbluth plot.** Here the weak contribution can also be separated from two-photon effects.

The cross-sections for right- and left-handed electrons and positrons have been evaluated\(^{13}\) in the Salam-Weinberg model assuming \(\sin^2 \theta_W = 0.25\). The cross-sections normalized to the pure one-photon cross-section are listed in Table II.9 as a function of \(x\) and \(y\), and are plotted in Fig. II.24 for \(x = 0.25\) as a function of \(y\). It is clear that there are large weak interaction effects in a kinematical range with acceptable counting rate (Fig. II.20). The helicity of the incident lepton beam must in general be known in order to study effects of the weak interaction. The transverse polarization of the circulating beam can be determined by measuring back-scattered circularly polarized laser light, as discussed in Chapter III, Section 4. This method has now been used successfully\(^{14}\) at SPEAR. The degree of longitudinal polarization at the interaction point can then be computed. It may also be possible to use \(\gamma + p \rightarrow \rho + p\) and \(\gamma + p \rightarrow \phi + p\) (s-channel helicity conservation) to monitor the electron helicity directly.
Table II.9
The cross-section $e + p \rightarrow e' + X$ for left- and right-handed electrons and positrons normalized to the one-photon exchange cross-section.
25 GeV electrons on 270 GeV protons.

<table>
<thead>
<tr>
<th>x/y</th>
<th>0.1</th>
<th>0.3</th>
<th>0.5</th>
<th>0.7</th>
<th>0.9</th>
<th>e_L^-</th>
<th>e_R^-</th>
<th>e_R^+</th>
<th>e_L^+</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.1</td>
<td>1.02</td>
<td>1.09</td>
<td>1.18</td>
<td>1.31</td>
<td>1.42</td>
<td>e_L^-</td>
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<td>0.68</td>
<td>0.53</td>
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<td>e_R^-</td>
<td>e_R^+</td>
<td>e_L^+</td>
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<td>1.02</td>
<td>1.16</td>
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<td>1.20</td>
<td>1.37</td>
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<td>e_L^+</td>
</tr>
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<td>e_L^-</td>
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<td>e_R^+</td>
<td>e_L^+</td>
</tr>
<tr>
<td>0.9</td>
<td>1.17</td>
<td>1.53</td>
<td>1.91</td>
<td>2.27</td>
<td>2.53</td>
<td>e_L^-</td>
<td>e_R^-</td>
<td>e_R^+</td>
<td>e_L^+</td>
</tr>
<tr>
<td></td>
<td>0.94</td>
<td>1.04</td>
<td>1.23</td>
<td>1.41</td>
<td>1.55</td>
<td>e_L^-</td>
<td>e_R^-</td>
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<tr>
<td></td>
<td>1.10</td>
<td>1.11</td>
<td>0.95</td>
<td>0.74</td>
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<td>e_L^-</td>
<td>e_R^-</td>
<td>e_R^+</td>
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<tr>
<td></td>
<td>0.89</td>
<td>0.72</td>
<td>0.54</td>
<td>0.36</td>
<td>0.24</td>
<td>e_L^-</td>
<td>e_R^-</td>
<td>e_R^+</td>
<td>e_L^+</td>
</tr>
</tbody>
</table>
3.4 Charged current events

High-energy electron-proton colliding rings make it possible to probe the charged weak current by observing the reaction $e + p \to \nu + X$. A measurement of $e + p \to \nu + X$, compared to present neutrino-induced reactions, offers several advantages:

i) A high, well-defined energy. The c.m. energy attainable with CHEEP corresponds to a monoenergetic neutrino beam of 14 to 20 TeV.

ii) Helicity. Both left- and right-handed electrons and positrons will be available, whereas the neutrinos have fixed helicity.

iii) Final state. The e-p collisions are equivalent to experiments with very low density targets, i.e. the produced particles escape from the interaction region and can be measured and identified. This is a great advantage compared to conventional neutrino experiments where the interaction usually takes place inside a calorimeter with poor identification power.

The charged current reaction $e + p \to \nu + X$, however, must be identified, and $Q^2$ and $\nu$ determined from a measurement of the produced hadrons only. This problem is similar to the one encountered in measuring the neutral current reaction $\nu + p \to \nu + X$. It will be discussed below in detail.

The number of charged current events produced per day in bins $dx dy = 0.04$ is plotted in Fig. II.25. These rates were evaluated$^{16}$ for a luminosity of $10^{32}$ cm$^{-2}$ sec$^{-1}$ assuming scaling and a point interaction, i.e. $m_W = m$. If this point-like behaviour would persist, then the weak interaction would become stronger than the electromagnetic interaction for $Q^2$ above 4000 GeV$^2$ -- a value well within the range of CHEEP -- and the rates would be acceptable for $Q^2$ up to 20,000 GeV$^2$.

However, in current models the weak interaction is damped by assigning a finite mass, of the order of 60 GeV, to the intermediate vector boson. This will reduce the rates by the characteristic propagator effect:

$$\left(\frac{m_W^2}{m_W^2 + Q^2}\right).$$

Therefore, unlike scaling violations, which according to present models are most striking at small values of $Q^2$, the existence of an intermediate vector boson will have little effect at small values of $Q^2$ but will change the cross-section by more than an order of magnitude at large $Q^2$. Furthermore, scaling violations are expected to affect the electromagnetic and the weak cross-section similarly, and can therefore be reduced by taking the ratio:

$$\frac{\sigma(e + p \to \nu + X)}{\sigma(e + p \to e' + X)}.$$

Next we discuss how charged current events can be recognized and separated from neutral current events, and then we describe two rather complementary methods which can be used to determine $Q^2$ and $\nu$. For this discussion we assume a detector such as the one described in Section 3.2 with FRAGCAL approximated as shown in Fig. II.26. The effect of the magnetic
field in FRAGCAL was neglected, as was the possibility of installing a calorimeter on the proton beam line beyond the proton quadrupoles. With such a calorimeter, angles down to a few mrad with respect to the proton direction can be covered.

As discussed in Section 3.1 a deep inelastic charged current event is characterized by a high-energy neutrino at large angles, with its transverse momentum balanced by the particles in the current jet and the proton jet. A neutral current event is obtained by replacing the neutrino with an electron.

We can now consider how to identify a charged current event. First, the transverse momentum of all hadrons -- neutral and charged -- with respect to the beam axis is measured, and events with a missing transverse momentum of more than 10 GeV are considered as candidates for deep inelastic events. As shown in Fig. II.25, such a cut does not substantially reduce the available kinematical region. It is, however, important to verify that the missing transverse momentum is not due to hadrons escaping down the aperture of FRAGCAL. To estimate this uncertainty, deep inelastic events were simulated as described in Section 3.1. Hadrons within one degree of the proton beam line impinging on FRAGCAL and were propagated through this detector, using a Monte Carlo shower program.

The uncertainty in $p_T^h$ caused by particles impinging on FRAGCAL is plotted in Fig. II.27 as a function of $x$ for $0.3 \leq y \leq 0.7$. This figure shows that the particles in the proton jet contribute less than 1 GeV/c to the uncertainty in $p_T^h$. The measured value of $p_T^h$ is determined by particles in the current jet which can be well measured. The cut on transverse momentum can therefore be used to exclude beam-gas or photoproduction events which will have a balanced transverse momentum. The events with $p_T^h > 10$ GeV are therefore largely due to deep inelastic processes, and the next step is to separate charged and neutral current events, i.e. to decide if a neutrino or an electron appeared opposite the current jet. The azimuthal angle is determined by a measurement of the transverse momentum; however, to determine the production angles unambiguously, the total hadron energy must also be known. The energy loss due to the hole in FRAGCAL and the amount of energy leaking through the sides of FRAGCAL have been computed as outlined above, and the results are shown in Fig. II.28. The leakage through the sides and the bottom of the calorimeter is rather small, and it can be further reduced by increasing the transverse dimensions of FRAGCAL. A substantial amount of energy is lost through the hole in the calorimeter, but note that this loss decreases with $x$ and accounts for only about 10% of the total laboratory energy for $x > 0.4$. However, the event-to-event fluctuations in the energy loss are rather large, as shown in Fig. II.29, and the production angle of the neutrino cannot be predicted precisely.

It is therefore important to cover nearly $4\pi$ in solid angle with no dead spots. The proposed detector satisfies these requirements. The ratio of charged to neutral current events is plotted in Fig. II.30 for $m_W = \infty$ and 63 GeV. The plot shows that the ratio of charged to neutral current events is nearly 1, i.e. an efficiency for electron detection of the order of 99% is equivalent to a 1% contamination of the charged current events with neutral current events. An efficiency of 99% seems realistic since the electron (Fig. II.9) appears at large angles with high energy and is therefore easily recognized. We therefore feel confident that the charged current events can be identified with only a small contamination of other events. This is sufficient to measure the total cross-section for $e + p \to \nu + X$. The rates for the total cross-section were evaluated for various assumptions.
and the results are plotted in Figs. 11.31 and 11.32. With CHEEP this cross-section can be explored for values of s between 4000 GeV$^2$ and 48,000 GeV$^2$. Note that the rates are sufficient even if the luminosity is reduced by a factor of 10. Within the present theories, left-handed positrons and right-handed electrons do not take part in the weak interaction. Hence a measurement of these cross-sections constitutes a very sensitive test of new weak currents.

More information can be obtained if the events are plotted as a function of $Q^2$ and $\nu$ or equivalently $x$ and $y$. These quantities can be determined from a measurement of the longitudinal and transverse momentum of the produced hadrons. We have seen above that the transverse momentum can be well measured, whereas a substantial amount of energy escapes along the proton direction. However, these losses are not very important at large values of $x$. The fraction of total hadronic energy in a deep inelastic event is plotted in Fig. 11.33 as a function of the production angle with respect to the proton direction for various values of $x$. It is possible, by installing detectors beyond the proton quadrupoles, to cover angles down to a few milliradians. It therefore seems feasible to determine $x$ and $y$ from a measurement of all hadrons for, say, $x \geq 0.4$. For smaller values of $x$, these quantities can be extracted from a measurement of the current jet only, using the following arguments.

The transverse momentum of the hadrons, which to a good approximation is given by the hadrons in the current jet, must be equal to the transverse momentum of the outgoing lepton:

$$ (p^L_T)^2 = (p^h_T)^2 = 4E_e E_p x y (1-y) . $$

This is a purely kinematical argument, and to proceed further we must assume that the hadrons are produced with limited transverse momentum and a reasonable distribution in rapidity, i.e. $[\exp (-p_T^2)]$, $[\exp (-\beta x)]$. Figures 11.15 and 11.16 show that a well-separated two-jet structure results in a wide range of values for $\alpha$ and $\beta$. The angles of the current jet can be expressed in terms of $x$ and $y$. Figure 11.34 shows lines of constant $x$ and $y$ as a function of $\theta_J$, the production angle of the current jet, and as a function of $(p_T^h)^2/4E_e E_p$. The errors corresponding to $\Delta \theta_J = \pm 3\%$ and $\Delta p_T/ p_T = \pm 0.1$ are shown. This method permits us to determine $x$ and $y$ with errors of less than 0.1 for $x \leq 0.5$ and $y \leq 0.7$. Note that these errors decrease rapidly with smaller values of $x$ and $y$. This method of determining $x$ and $y$ would clearly be unsatisfactory if the transverse momentum distribution of the produced hadrons became rather flat. Such an observation would be in direct contradiction to all our prejudices and would constitute a major discovery by itself.

In summary we have two complementary methods of determining $x$ and $y$ for a charged current event. The first method relies upon measuring the angles and the momenta of all hadrons in the final state; the second method derives the quantities from a measurement of the current jet only. Both methods can be tested and the errors evaluated using neutral current events.

Plotting the charged current cross-section as a function of $x$ and $y$ makes it possible to determine the propagator effects with increased sensitivity. This is shown in Figs. 11.30 and 11.25. Figure 11.35 shows that the cross-section for $0.4 > x > 0.6$ changes by more than an order of magnitude when the point-like coupling is replaced with $m_W = 63$ GeV.
4. Final States with Neutral and Charged Currents

Detailed studies of the final state require rather sophisticated detectors capable of identifying and measuring charged and neutral particles over a wide momentum range. To quantify these statements, we consider two important final-state experiments in more detail: the search for new flavours, and the search for new electron-like leptons.

4.1 New flavour production

Electroproduction is expected to be a fertile source of hadrons with new flavours. A specific production mechanism analogous to electromagnetic pair production in the field of a nucleus is shown in Fig. II.36. The sea consists of $qar{q}$ pairs, and a photon with $Q^2 > m_q^2$ is able to resolve this structure. The large $Q^2$ photon interacts with one of the quarks and displaces it to a distant region of phase space where it has a low probability of recombining with its partner. In this case the quark will more likely recombine with one of the old quarks and appear as a hadron with a new flavour in the current jet. Its partner will most likely show up as a new hadron in the proton fragmentation jet. The rate for this process is listed in Table II.10 (taken from Ref. 10).

The kinematics are favourable for detecting the new hadron in the current jet. The $qar{q}$ pairs are concentrated at low $x$, i.e. the new hadron will appear at large angles and the multiplicity and the total momentum of the current jet will be rather low. This is shown in Fig. II.37, where the momentum of the current fragment versus $y$ is plotted. Also shown is the kinematic boundary for $\theta_J = 30^\circ$ and $\theta_J < 150^\circ$. The signal-to-background ratio seems to be favourable -- a new quark with a mass around 10 GeV and charge $2/3e$ will contribute of the order of 5% of the total cross-section at small $x$.

Table II.10

Number of $e + p \rightarrow e + X$ events/day (for $L = 10^{32}$ cm$^{-2}$ sec$^{-1}$) in which scattering occurs on a quark with a new flavour (leading to associated production) for quark charge $2/3$, mass $M_t$ and $Q^2 > Q_{min}$ (according to assumptions discussed in the text). The numbers in brackets ( ) are the numbers of events off the old (u, d, s, and c) quarks. $Q^2$ and $s$ are in units of GeV$^2$.

<table>
<thead>
<tr>
<th>$Q^2_{min}$</th>
<th>$s = 6720$</th>
<th>$s = 27,000$</th>
<th>$s = 40,000$</th>
</tr>
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<tr>
<td>$M_t = 5$ GeV</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>25</td>
<td>5,500 (10$^5$)</td>
<td>15,000 (10$^5$)</td>
<td>18,400 (10$^5$)</td>
</tr>
<tr>
<td>100</td>
<td>550 (10$^4$)</td>
<td>3,400 (2 x 10$^4$)</td>
<td>4,600 (2 x 10$^4$)</td>
</tr>
<tr>
<td>200</td>
<td>180 (4000)</td>
<td>1,200</td>
<td>1,600</td>
</tr>
<tr>
<td>$M_t = 10$ GeV</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>100</td>
<td>370 (10$^4$)</td>
<td>1,600 (8 x 10$^4$)</td>
<td>2,200 (8 x 10$^4$)</td>
</tr>
<tr>
<td>400</td>
<td>17 (900)</td>
<td>220 (6000)</td>
<td>350 (6000)</td>
</tr>
<tr>
<td>800</td>
<td>2.3 (200)</td>
<td>45 (900)</td>
<td>82 (900)</td>
</tr>
<tr>
<td>$M_t = 25$ GeV</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>625</td>
<td>1.7 (400)</td>
<td>33 (1600)</td>
<td>57 (1600)</td>
</tr>
<tr>
<td>2,500</td>
<td>- (2)</td>
<td>1 (70)</td>
<td>2.8 (130)</td>
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<tr>
<td>5,000</td>
<td>-</td>
<td>0.06 (8)</td>
<td>0.3 (20)</td>
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<tr>
<td>2,500</td>
<td>-</td>
<td>0.2 (70)</td>
<td>0.4 (130)</td>
</tr>
<tr>
<td>10,000</td>
<td>-</td>
<td>-</td>
<td>-</td>
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</table>
The new hadrons can be identified by their semileptonic decays: \( H \to e^0 X \to (e + \bar{u} + X) \), i.e. \( e + p \to e' + X \), \( p \to \bar{u} + X \). It is unlikely that this signature is sufficient to pin down a certain flavour since presumably there are several sources of inclusive electrons in the current jet. However, the presence of single electrons (or muons), or events with both electrons and muons, in the current jet are a clear signature for new physics. The new hadrons will appear at large angles with moderate momenta (Fig. II.37), and it might therefore be possible to identify the decay \( H \to \) hadrons. Provided there is neutral and charged particle identification, this seems entirely feasible, particularly since we expect a rather favourable signal-to-background ratio.

New flavours can also be produced by the charged current, either off the valence quarks or off the sea quarks. The new hadron is produced singly and it travels with the particles in the current jet. In this case, masses up to 50-75 GeV may be probed, but the background may be more severe and particles of higher momenta must be identified.

Experiments to search for new flavours require a magnetic detector with good particle identification and also the ability to reconstruct \( \pi^0 \)'s and \( \eta \)'s.

4.2 New lepton production

New leptons with electron quantum numbers can be produced in electron-proton collisions via the diagrams depicted in Fig. II.38a. The rates for \( s \gg m_E^2 \) are related to the rates for the usual charged current events as follows:

\[
\frac{\sigma(e_1^+ + p \to E^0 + X)}{\sigma(e_1^- + p \to \nu + X)} = \frac{\left(\frac{\sigma_{e1e1E}}{\sigma_{e1e1\nu}}\right)^2}{\frac{\sigma(e_1^- + p \to \bar{E}^0 + X)}{\sigma(e_1^- + p \to \bar{\nu} + X)} = \frac{\sigma(e_1^+ + p \to E^0 + X)}{\sigma(e_1^+ + p \to E^+ + X)} = \frac{\sigma_{e1e1E}}{\sigma_{e1e1\nu}}}
\]

\[
\frac{\sigma(e_2^- + p \to E^0 + X)}{\sigma(e_2^- + p \to \bar{E}^0 + X)} = \frac{\left(\frac{\sigma_{e2e2E}}{\sigma_{e2e2\nu}}\right)^2}{\frac{\sigma(e_2^+ + p \to E^0 + X)}{\sigma(e_2^+ + p \to E^+ + X)} = \frac{\sigma_{e2e2E}}{\sigma_{e2e2\nu}}}
\]

\[
\frac{\sigma(e_3^- + p \to E^0 + X)}{\sigma(\nu + p \to \nu + X)} = \frac{\sigma(e_3^+ + p \to E^0 + X)}{\sigma(\nu + p \to \nu + X)} = \frac{\sigma_{e3e3E}}{\sigma_{e3e3\nu}}
\]

These relationships were taken from Ref. 10, and the approximate equalities are valid in a valence parton model. The asymptotic rates are approached fairly rapidly, as shown in Fig. II.38b taken from Ref. 17.
Using this curve and the total cross-section rates plotted in Figs. II.31 and II.32, valid for $m_W = 63$ GeV and asymptotic freedom, we obtain the production rates for a heavy lepton assuming:

$$g_{e_LE^0} = g_{e_LV}.$$  

This rate is listed in Table II.11 as a function of lepton mass for 25 electrons on 270 GeV protons or 30 GeV electrons on 400 GeV protons and a luminosity of $10^{32}$ cm$^{-2}$ sec$^{-1}$. Since the leptons produced have electron quantum numbers, they might have rather distinct, neutrinoless decays. Some of the anticipated decay modes are:

$$
egin{align*}
E^0 &\to e^- + \text{hadrons} & E^- &\to e^- + \text{hadrons} \\
\to e^- + \mu^+ + \nu_\mu &\quad & \to e^- + e^+ + e^- \\
\quad &\to \nu_\mu + \text{hadrons} & \to \bar{e}^ - + \mu^- + \bar{\nu}_\mu \\
\quad &\to \nu_\mu + e^+ + e^- & \to \nu_\mu + \bar{\nu} \\
\quad &\to \nu_\mu + \mu^+ + \mu^- & \to \nu_e + \text{hadrons} \\
\quad &\to E^- + \text{hadrons} & \to E^0 + \text{hadrons} \\
\quad &\to E^- + \mu^+ + \nu_\mu & \to E^0 + \mu^- + \nu_\mu \\
\quad &\to E^- + e^+ + \nu_e & \to E^0 + e^- + \nu_e
\end{align*}
$$

<table>
<thead>
<tr>
<th>Table II.11</th>
</tr>
</thead>
<tbody>
<tr>
<td>Production rate for new leptons in $e_L^- + p \to E^0 + X$.</td>
</tr>
<tr>
<td>Luminosity $10^{32}$ cm$^{-2}$ sec$^{-1}$, asymptotic freedom and $m_W = 63$ GeV assumed.</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Mass of lepton (GeV)</th>
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<th>50</th>
<th>75</th>
<th>100</th>
<th>125</th>
</tr>
</thead>
<tbody>
<tr>
<td>Events/day ($s = 27,000$ GeV$^2$)</td>
<td>419</td>
<td>176</td>
<td>62</td>
<td>15</td>
<td>-</td>
</tr>
<tr>
<td>Events/day ($s = 48,000$ GeV$^2$)</td>
<td>620</td>
<td>361</td>
<td>199</td>
<td>87</td>
<td>31</td>
</tr>
</tbody>
</table>

How can these leptons be detected? The kinematics for heavy lepton production is expected to be similar to the kinematics for deep inelastic electroproduction, i.e. we expect to observe a three-jet structure (Fig. II.38a) with the decay products from the lepton jet well separated from the proton and the current jet. The existence of a neutral lepton can be inferred from the observation of such events with one jet consisting of two charged leptons, i.e. $(e^- \mu^+)$, $(e^- e^+)$, or $(\mu^- \mu^+)$. The mass of a neutral lepton can be determined directly from the decay to $e^- +$ hadrons. The existence of a new charged lepton can be inferred and its mass determined by observing jets consisting of three charged leptons $(e^- e^+ e^-)$ or $(e^- \mu^+ \mu^-)$.  

The decay products from a heavy lepton of mass less than 50 GeV are in general well separated from the particles belonging to the current jet. This is shown in Fig. II.39 for the decay mode $E^+\nu^+e^+\bar{e}^+$, $E^-\nu^-e^-\bar{e}^+$, i.e. the decay into three massless particles. The plot is made for 25 GeV electrons on 270 GeV protons producing a heavy lepton of mass 20 GeV. The number of tracks is plotted as a function of the angle between the axis of the current jet and particles from the decay of the new lepton. The current jet has a width of a few degrees and is clearly separated from the decay products of the E. The leptons from the decay have an average momentum around 15 GeV, as shown in Fig. II.40, and are therefore rather easy to identify and measure using calorimeters (or magnets for the low-energy ones).

The same plots are shown in Figs. II.41 and II.42 for a heavy lepton of mass 75 GeV. The topology resulting from the production of such a heavy object is markedly different from that of a conventional deep inelastic event and should be easy to recognize. It will be possible to separate the current jet from the decay products of the heavy leptons if the particles belonging to the current jet are within a cone with an opening angle of a few degrees as expected (Table II.7).

The heavy leptons thus have quite spectacular signatures, and it seems feasible to identify leptons with a production rate of the order of a few events per day. Thus with a luminosity of $10^{32}$ cm$^{-2}$ sec$^{-1}$, it is possible to explore the mass range up to 100 GeV for $s = 27,000$ GeV$^2$ or up to 150 GeV for $s = 48,000$ GeV$^2$.

4.3 Detector

A general detector suitable for investigating the final state in deep inelastic reactions can be obtained from the non-magnetic detector, discussed in Section 3.2, by replacing the central unit (CENCAL) with a magnetic detector and retaining the calorimeter elements which cover angles close to the beam line. This central magnetic detector must fulfil the criteria discussed below:

1) Provide magnetic analysis over a large solid angle. A solenoid with the axis along the beam direction can produce a uniform field over a large solid angle, and it is a natural choice for the magnetic analysis of particles produced at large angles with respect to the beam axis. The c.m. system in an e-p ring is not at rest but moves along the proton direction, and it might therefore be advantageous to centre the solenoid downstream from the interaction point. A toroidal field, produced by thin superconducting coils, can be used to analyse particles down to roughly 15° with respect to the proton beam line. It might even be possible to combine these two detectors. However, it is most important for the measurement of charged current events that the detector has no dead regions.

2) Identify hadrons which travel within a jet over a wide range in momentum. Several new schemes have been proposed and will be tried out at the detectors planned for PETRA and PEP. It seems possible to replace high-pressure gas counters with counters using aerogel as the radiator. Samples of aerogel with a refractive index between 1.008 and 1.1, corresponding to a threshold for pions (kaons) of 1.18 GeV/c (4.29 GeV/c) and 0.33 GeV/c (1.20 GeV/c), respectively, have been produced$^{18}$. Ionization measurements$^{19}$ or imaging $\gamma$ Cerenkov counters$^{20}$ can also be used to identify the hadrons. In the former case, particles with a known momentum can be identified by measuring the energy loss a few hundred times along the track.
3) Identify and measure electrons within the current jet. This can be achieved using finely grained liquid-argon calorimeters mounted outside the magnetic field region. Such calorimeters are now being built for PEP and PETRA. A finely grained calorimeter is also needed to measure \( \eta \)'s and \( \eta \)'s over a large momentum range. This is an important requirement if we want to observe decay modes like \( H \to \text{hadrons} \) or \( L \to e^+ \text{hadrons} \).

4) Detect muons over a limited solid angle. This can be obtained by mounting track chambers after the magnetic flux return yokes.

Many of the features needed for a large general detector that would be used to study electron-proton collisions are already envisaged for the large detection devices proposed for PEP and PETRA. It therefore seems unrealistic to attempt a detailed layout of a detector to study final states in \( e-p \) collisions before the experience gained with these devices has been evaluated.

5. PHOTOPRODUCTION

5.1 The beam

Electroproduction in the limit of \( Q^2 \to 0 \) can be described as the radiation of an almost real photon followed by the interaction of the photon with the proton. The electron beam is therefore equivalent to a well-collimated bremsstrahlung beam with equivalent end point energies of 14.4 TeV (25 GeV electrons on 270 GeV protons) or 25.6 TeV (30 GeV electrons on 400 GeV protons). The intensity of the beam, i.e. the number of photons with energies between \( k \) and \( k + dk \) per incident electron multiplied by the photon energy \( k \), is given by the well-known Weizsäcker-Williams formula:

\[
k \cdot N(k) \cdot dk = \frac{\alpha}{2\pi} \left( \frac{E^2 + E'^2}{E^2} \right) \ln \left( \frac{Q_{\text{max}}^2}{Q_{\text{min}}^2} \right) dk,
\]

where \( E \) and \( E' \) are the energies of the incident and the scattered electron, and \( Q^2 = EE' \cdot q^2 \), \( k = E - E' \). The equivalent photon energy \( \nu \) in the proton rest system is given by: \( \nu = 2k \cdot E_p/m_p \)

where \( E_p \) is the energy of the proton. Not tagging the scattered electrons, yields the highest luminosity at the price of leaving the photon energy unknown. This is acceptable for a whole series of experiments in which the photon energy can be determined from a measurement of the final state. The minimum momentum transfer squared for an untagged photon beam is given by:

\[
Q_{\text{min}}^2 = \frac{m_e^2 (E - E')^2}{EE'}.
\]

The intensity \( k \cdot N(k) \) is plotted in Fig. II.43. The effective luminosity for photon-proton collisions \( L(\gamma-p) \) is approximately 0.05 \( L(e-p) \) for photons with energies above those which can be reached at the SPS. The photons will be circularly polarized if the incident electron beam has a well-defined helicity.

The energy of the scattered photon can be determined from a measurement of the angle and the energy of the scattered electron. Assuming that the electrons are tagged using a liquid-argon calorimeter with \( dE'/E' = \pm 0.05/\sqrt{E'} \) GeV, we find \( d\nu/\nu = \pm 0.045 \) for \( \nu = 2879 \) GeV and \( d\nu/\nu = \pm 0.0021 \) for \( \nu = 13,817 \) GeV. The angular range to be covered by the tagging counters
is limited at small angles by the minimum distance allowed between the counter and the beam, and by the length of the interaction region. We assume $\theta_{\text{min}}$ to be 5 mrad, corresponding to 2.5 cm at 5 m. To be able to neglect the longitudinal part of the photon, we rather arbitrarily assume $Q^2_{\text{max}} \leq 0.02$ GeV$^2$. The angular range to be covered by the tagging counters for a 25 GeV electron beam is shown in Fig. II.44. Note that "low-energy $\gamma$'s", i.e. $\nu \leq 5000$ GeV, have only a small tagging efficiency. This is reflected in the photon intensity spectrum also plotted in Fig. II.45. The "low-energy region" is better explored by reducing the incident electron energy and correspondingly increasing the angles. The tagging system should cover angles between 5 mrad (or less) and 100 mrad, and should measure the energy and the direction of the scattered electron. Note that such a measurement determines the polarization of the photon with respect to the production plane defined by the incident photon and some outgoing hadron.

5.2 Experiments

With an e-p luminosity of $10^{32}$ cm$^{-2}$ sec$^{-1}$, of the order of $0.4 \times 10^8$ events/day with a $Q^2$ less than $m_e^2$ are expected using an untagged beam, or approximately $3 \times 10^8$ events/day using the tagged beam described above. There is therefore a wide range of experiments that can be investigated with a high-energy photon beam, as discussed elsewhere in this report. Here we concentrate on three typical experiments: Compton scattering, diffractive production of $l^-$ states, and Primakoff production of resonances with positive charge conjugation and spin different from one. The Feynman graphs for the last two processes are shown in Figs. II.45 and II.46.

5.2.1 Compton scattering

Compton scattering is one of the most basic processes that can be investigated with an incident photon beam, and a measurement of this reaction might reveal properties that are unique to the photon. The reaction has very simple kinematics: the scattered photon travels forward (along the electron direction) with approximately the incident energy, and the invariant momentum transfer squared $t$ is given by $t = 4kk' \sin^2 \theta/2 = k^2 \bar{e}_y^2$. The proton continues with the incident energy at a slightly different angle $\theta_\text{p} = \sqrt{t/E_\text{p}}$.

At small values of $t$ (say, $t \leq 0.5$ GeV$^2$) only the scattered electron and the scattered photon can be detected in coincidence, and this can be done using the finely grained forward-tagging detector. Note that the scattered electron and photon have to be in a plane which includes the beam axis -- this, together with the requirement that no other particles should be observed, will hopefully lead to a clean trigger. With the calorimeter covering angles from 5 mrad to 300 mrad we can cover the $t$-range from $2.5 \times 10^{-5}$ GeV$^2$ to 0.09 GeV$^2$ for $k = 1$ GeV ($\nu = 570$ GeV), and from $t = 0.014$ GeV$^2$ to 52 GeV$^2$ for $k = 24$ GeV ($\nu = 13,817$ GeV). At larger $t$-values the recoil proton will travel at angles which are large compared to the beam divergence ($\theta \sim \sqrt{E/E_\text{p}} \sim \sqrt{0.5/270} = 2.6$ mrad) and may be measured. It might be sufficient to measure only the recoil proton in coincidence with the scattered photon; this would increase the yield by an order of magnitude.

Extrapolating the present data on Compton scattering to higher energies would give a yield of 3000 tagged events/day integrated over all $t$. If it would suffice to measure only the scattered photon and the recoil proton in coincidence, the rate of events with $|t| > 0.5$ would increase from 120 events/day to about 1000 events/day.
5.2.2 Production of vector mesons

Vector mesons have the same quantum number as photons and therefore we might expect them to be diffractively produced by an incident photon beam. This is expressed in the vector dominance model where the photoproduction cross-section is proportional to the cross-section of elastic scattering between vector meson and proton as shown in Fig. II.45 (taken from Ref. 10). Photoproduction of vector mesons can therefore be used to obtain information on elastic meson-proton scattering at high energies. These experiments can be made very sensitive by using polarized photons and by determining the helicity of the produced vector meson. The rates for the old vector mesons are listed in Table II.12 (from Ref. 10).

Table II.12

<table>
<thead>
<tr>
<th>Process</th>
<th>$\sigma_{\text{tot}}$</th>
<th>Events/day not tagged</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\gamma + p \rightarrow p + p$</td>
<td>13 $\mu$b</td>
<td>$6 \times 10^6$</td>
</tr>
<tr>
<td>$\gamma + p \rightarrow \omega + p$</td>
<td>1.4 $\mu$b</td>
<td>$0.6 \times 10^6$</td>
</tr>
<tr>
<td>$\gamma + p \rightarrow \phi + p$</td>
<td>0.5 $\mu$b</td>
<td>$2.2 \times 10^5$</td>
</tr>
<tr>
<td>$\gamma + p \rightarrow J/\psi + p$</td>
<td>30 nb</td>
<td>$1.5 \times 10^6$</td>
</tr>
</tbody>
</table>

The energy and the momentum of a vector meson produced at $p_T = 0$ by a photon of energy $k$ is to a good approximation given by

$$E_V = \frac{4k^2 + m_V^2}{4k}$$

$$p_V = \frac{4k^2 - m_V^2}{4k},$$

where $m_V$ is the mass of the vector meson. The produced vector meson will therefore travel along the electron direction if $2k > m_V$, or along the direction of the proton if $m_V > 2k$.

The light vector mesons will therefore travel along the electron direction with approximately the same energy as that of the incident photon. The $\rho$ and the $\phi$ can be identified using the characteristic decay modes $\rho \rightarrow \pi^+ + \pi^-$ and $\phi \rightarrow k^+ + k^-$. The opening angle between the $\pi^+$ and the $\pi^-$ from $\rho \rightarrow \pi^+ + \pi^-$ and their energies are plotted in Fig. II.47 as a function of $\cos \theta$ ($\theta$ is the decay angle in the rest system of $V$ for various values of $k$). The energies and the opening angles are such that these particles can be measured. Using a tagged photon beam, it seems sufficient to measure the opening angle and, roughly, the energies in order to identify $\rho \rightarrow \pi^+ + \pi^-$ and $\phi \rightarrow k^+ + k^-$. A photon beam is expected to be a copious source of new vector mesons. The maximum observable mass is determined by the minimum momentum transfer to the proton:

$$t_{\text{min}} = \frac{m_V^2 - m_p^2}{s^2};$$
\( t_{\text{min}} = 0.2 \) corresponds to a vector meson of mass 113 GeV. However, the cross-section for producing vector mesons with new flavours is expected to decrease with increasing mass as \( 1/M_V^2 \). It might thus be possible to search for new \( V^- \) resonances with a mass up to maybe 30 GeV. These particles have a very clean signature: most of the decays will yield either lepton pairs or two hadron jets, well separated in space. The decay products will appear at large angles where they can be easily detected. In Fig. II.48 the opening angle and the momentum of the decay electron for \( V^- \rightarrow e^+ e^- \) are plotted versus \( \cos \theta^* \) for \( m_V = 30 \) GeV and different photon energies.

5.2.3 Primakoff effect

An incident photon can interact with the Coulomb field of the proton to produce hadrons with spin different from one and with positive charge conjugation. This process, shown in Fig. II.46, leads to a differential cross-section of the form\(^{22}\)

\[
\frac{d\sigma}{dt} = \frac{8\pi \alpha}{m^2} \left( \Gamma_{\gamma\gamma}/m \right) (t - t_{\text{min}})/t^2 |F(t)|^2.
\]

In this formula, \( m \) is the mass of the produced particle, \( \Gamma_{\gamma\gamma} \) the partial width of the resonance for decay to a pair of photons, \( t_{\text{min}} \) the minimum momentum transfer to the proton \( [t_{\text{min}} = (m^2/2k)^2] \), and \( F \) the photon form factor. Since the proton form factor is near to one, this yields a total cross-section:

\[
\sigma = (8\pi \alpha/m^2)(\Gamma_{\gamma\gamma}/m)(\ln t_{\text{max}}/t_{\text{min}} + 1),
\]

which increases logarithmically with energy. The rates for some particles with a large value of \( \Gamma_{\gamma\gamma} \) are listed in Table II.13.

<table>
<thead>
<tr>
<th>Particle</th>
<th>( \Gamma_{\gamma\gamma} ) (MeV)</th>
<th>( \sigma ) (nb)</th>
<th>Untagged events/day</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \pi^0 )</td>
<td>( 8 \times 10^{-6} )</td>
<td>5.64</td>
<td>2436</td>
</tr>
<tr>
<td>( \eta )</td>
<td>( 3.8 \times 10^{-4} )</td>
<td>3.06</td>
<td>1322</td>
</tr>
<tr>
<td>( \chi^0 )</td>
<td>( 2 \times 10^{-4} )</td>
<td>0.27</td>
<td>115</td>
</tr>
<tr>
<td>( \xi )</td>
<td>( \sim 4 \times 10^{-3} )</td>
<td>2.13</td>
<td>920</td>
</tr>
<tr>
<td>( \eta_c )</td>
<td>( \sim 20 \times 10^{-3} )</td>
<td>0.76</td>
<td>328</td>
</tr>
</tbody>
</table>

This process makes it possible to determine radiative widths \( \Gamma_{\gamma\gamma} \) directly for a series of resonances. The rates are also large enough to measure the cross-section as a function of \( q^2 \). This is related to the structure of the particle. The reaction is also well suited to searching for states with new flavours. Figure II.49 shows the kinematics for \( \gamma p \rightarrow Xp \rightarrow \gamma \gamma p \) for \( M_X = 10 \) GeV and various photon energies.
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22) H. Primakoff, Phys. Rev. 81 (1951) 899.
Fig. II.1 Front face of quadrupole Q3

Fig. II.2 Experimental hall considered for the pp project
Fig. II.3 Synchrotron radiation spectra from the magnets near the interaction region
Fig. II.4 Machine elements upstream of the septum magnet S2

Fig. II.5 Horizontal layout of the vacuum chamber and the synchrotron radiation absorbers

Fig. II.6 Vertical layout of the vacuum chamber
Fig. II.7 Time distribution of triple coincidences in the "far" telescope for three consecutive SPS cycles

Fig. II.8 Plots of beam loss rate and telescope counting rate for the SPS coasting beam as a function of storage time
Fig. II.9  Feynman graph for deep inelastic e-p reactions

Fig. II.10  Kinematical boundaries in the $\nu$, $Q^2$ plane for 25 GeV electrons on 270 GeV protons. (Lines of constant energy and angle of the scattered electron are dotted. The solid lines represent fixed mass of the produced hadron system.)
Electron Proton Collision at 25 and 270 GeV

Polar diagram of the outgoing lepton
Lines of constant \( W \) and \( Q^2 \)

Fig. II.11 Polar diagram of the scattered electron. The ellipses and the parabolas represent constant \( W \) and \( Q^2 \).

Event viewed along the beam direction

Fig. II.12 Event topology expected in a deep inelastic \( e-p \) interaction
Electron Proton Collision at 25. and 270. GeV

Polar diagram of the fragments of the current jet
Lines of constant $W$ and $Q^2$

Fig. II.13 Polar diagram of the fragments of the current jet. The lines represent constant $W$ and $Q^2$.

Electron Proton Collision at 25. and 270. GeV

Polar diagram of the outgoing lepton, with fragments directions
Lines of constant $W$ and $Q^2$

Fig. II.14 Polar diagram of the scattered lepton. The direction and the momentum (divided by 10) of the current and of the proton fragments are shown. The lines represent constant $W$ and $Q^2$. 
Fig. II.15  Distributions of the produced hadrons in the production angle $\theta$ for $(x, y) = (0.3, 0.9)$ and $(0.9, 0.3)$. The effect of altering the rapidity distribution of the hadrons is shown.

Fig. II.16  Distribution of the produced hadrons in the production angle for $(x, y) = (0.5, 0.5)$. The effect of changing the $p_T$ distribution is shown.
Fig. II.17 Deep inelastic muon scattering event (Ref. 6)

Fig. II.18 A compact non-magnetic detector for inclusive measurements of neutral and charged current events
Fig. II.19 The non-magnetic detector viewed along the beam direction. The boundary of the SPS tunnel is shown.

Fig. II.20 The number of electroproduction events per day in bins of $\Delta x = 0.04$ as a function of $Q^2$ and $\nu$. The rates are evaluated for one-photon exchange assuming scaling and a luminosity of $10^{32}$ cm$^{-2}$ sec$^{-1}$. These rates are changed in an asymptotically free theory by the factor listed in the brackets.
**Fig. II.21** The number of electroproduction events per day with $Q^2 > Q_0^2$ as a function of $Q_0^2$. The rates were evaluated for one-photon exchange assuming scaling and a luminosity of $10^{32} \text{ cm}^{-2} \text{ sec}^{-1}$.

**Fig. II.22** The kinematical region in a $Q^2, \nu$ plot which can be populated by beam-gas events.
**Fig. II.23** The predicted resolution in $d\psi/\psi$ and $dQ^2/c^2$ for a measurement of $e + p \rightarrow e' + X$ using CENCAL.

**Fig. II.24** The cross-section $e + p \rightarrow e' + X$ for left- and right-handed electrons and positrons normalized to the one-photon exchange cross-section.
Fig. II.25 The number of events in bins of $dxdy = 0.04$ predicted for the reaction $e^- p \rightarrow \nu + X$ as a function of $Q^2$ and $\nu$. The rates are evaluated for $s = 27,000$ GeV$^2$ assuming scaling and a point-like coupling.

Fig. II.26 Idealized version of FRAGCAL used in the Monte Carlo computations.
Fig. II.27 Transverse momentum of particles from the reaction $e + p \rightarrow \nu + X$ in FRAGCAL as a function of $x$.

Fig. II.28 Energy loss in FRAGCAL as a function of $x$ evaluated for the reaction $e + p \rightarrow \nu + X$.

Fig. II.29 Fluctuations of the energy loss in FRAGCAL.
Fig. II.30 The ratio \((e^+_\mu + p \rightarrow \nu + X)/(e^+_\mu + p \rightarrow e' + X)\) as a function of \(\nu\) and \(Q^2\) for \(s = 27,000\ \text{GeV}^2\). The ratio is evaluated assuming scaling, a point-like weak coupling, and \(m_\mu = 63\ \text{GeV}\).

Fig. II.31 The number of events per day for \(e^+_\mu + p \rightarrow \nu + X\) as a function of energy.
**Fig. II.32** The number of events per day for \( e_R^* + p \to \bar{\nu} + X \) as a function of energy

**Fig. II.33** The fraction of the total hadronic energy predicted for the reaction \( e^- + p \to \nu + X \) in bins of 2.5 mrad, as a function of the opening angle with respect to the proton direction
**Fig. 11.34** Production angle $\theta$ of the current jet versus $p_T^2/4E_pE_p$. Lines of constant $x$ and $y$ are shown. The uncertainty in $x$ and $y$ resulting from a measurement of $p_T$ and $\theta_j$ is indicated by rectangles.

**Fig. 11.35** Number of events per 100 days for the reaction $e^+_L + p \rightarrow \nu + X$ as a function of $y$ for $0.4 < x < 0.4$ and $0.6 < x < 0.6$. The rates are evaluated assuming a luminosity of $10^{32}$ cm$^{-2}$ sec$^{-1}$, asymptotic freedom, and $m_w = \infty$ and 63 GeV.
Fig. II.36 Mechanism for associated electroproduction of new flavours.

Fig. II.37 Momentum of the current fragment as a function of $y$ for low values of $x$.

Fig. II.38 a) Mechanism and topology for heavy lepton production in $e-p$ collisions; b) Predicted production rate of heavy leptons normalized to the asymptotic value as a function of $s/M_{E}^{2}$ (Ref. 17).
Fig. II.39 Number of electrons from $E^+ \rightarrow e^- + e^+ + e^-$ versus the angle $\theta$ between the electron and the current jet; evaluated for $s = 27,000 \text{ GeV}^2$ and a lepton of $20 \text{ GeV}$ mass.

Fig. II.40 Number of electrons from $E^+ \rightarrow e^- + e^+ + e^-$ versus momentum; evaluated for $s = 27,000 \text{ GeV}^2$ and a lepton of $20 \text{ GeV}$ mass.

Fig. II.41 Number of electrons from $E^- \rightarrow e^+ + e^- + e^-$ versus $\phi$; evaluated for $s = 27,000 \text{ GeV}^2$ and a lepton of $75 \text{ GeV}$ mass.

Fig. II.42 Number of electrons from $E^- \rightarrow e^- + e^+ + e^-$ versus momentum; evaluated for $s = 27,000 \text{ GeV}^2$ and a lepton of $75 \text{ GeV}$ mass.
Fig. II.43 Intensity of the virtual photon spectrum versus photon energy for a tagged ($Q_\text{max}^2 = 0.02$ GeV$^2$) and an untagged electron beam.

Fig. II.44 Angular range to be covered by the tagging system for a 25 GeV incident electron beam.

Fig. II.45 Feynman graph for diffractive production of vector mesons.

Fig. II.46 Feynman graph for Primakoff production.
Fig. II.47 Kinematics for $\gamma + p \rightarrow \rho + p \rightarrow \pi^+\pi^- + p$ produced at $p_T = 0$. Plotted are a) the $\pi^+\pi^-$ opening angle and b) the pion momentum in the machine system versus $\cos \theta^*$ for different photon energies; $\theta^*$ is the decay angle in the $\rho$ rest system.

Fig. II.48 Kinematics for $\gamma + p \rightarrow V + p \rightarrow e^+e^- + p$ produced at $p_T = 0$. The mass of $V$ is 30 GeV. Plotted are a) the angle of the leptons and b) the lepton momentum in the machine system versus $\cos \theta^*$ for different photon energies.

Fig. II.49 Kinematics for $\gamma + p \rightarrow X + p \rightarrow \gamma + \gamma + p$ produced at $p_T = 0$. The mass of the produced meson is 10 GeV. Plotted are a) the $\gamma\gamma$ opening angle and b) the photon energy in the machine system versus $\cos \theta^*$ for different incident photon energies.
CHAPTER III

AN e-p COLLIDING BEAM FACILITY WITH THE CERN SPS

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1. **INTRODUCTION**

The possibility of an e-p colliding beam machine was first discussed by Hereward et al.\textsuperscript{1) and further examined by Goldzah\l and Michaelis\textsuperscript{3) in connection with the CERN Intersecting Storage Rings (ISR). In 1971 the interest in e-p collisions was revived by an LBL/SLAC group\textsuperscript{4). An e-p option has also been considered for ISABELLE\textsuperscript{5), in DESY\textsuperscript{6), for the EPIC facility\textsuperscript{6), for POPAC\textsuperscript{7), and for the CERN Large Storage Rings (LSR)\textsuperscript{8). The possibility of colliding the protons of the FNAL main ring with electrons stored in a separate storage ring housed in the same tunnel was examined during an FNAL summer study in 1976\textsuperscript{9).}

Colliding the 400 GeV Super Proton Synchrotron (SPS) proton beam, circulating in a bypass, with an electron beam of 5 GeV, stored in a separate small ring, was considered previously\textsuperscript{10).}

In November 1976, B.H. Wilk proposed that CERN investigate whether e-p collisions could be obtained at the SPS by adding an electron storage ring to the proton synchrotron. A preliminary study\textsuperscript{11) showed that such an addition would provide a reasonable luminosity, and a study group was therefore formed in February 1977 to examine the design problems in more detail.

In October 1977 a study week, sponsored by ECPA, was held at Rutherford Laboratory where a status report on CHEEP was presented, as well as an updated version of the proposal for an e-p facility at DESY\textsuperscript{12). The concept of CHEEP was scrutinized critically at this meeting, and valuable advice and encouragement was obtained.

Work on this proposal, the CERN High-Energy Electron Proton (CHEEP) Facility, was first directed towards establishing the basic feasibility of the project. It soon became apparent that assessing the impact of such a facility on the SPS programme required detailed design work extending beyond the limited time and effort immediately available. We have therefore concentrated on a number of points which appear to be of crucial importance, and the present status report indicates the progress to date. Continuation of this work is needed to bring the studies up to the level required for a full design report.

In the scheme advanced here, an electron storage ring is housed in the SPS tunnel above the synchrotron, permitting a relatively high energy at modest RF power owing to the large radius of the SPS. The nominal electron energy is 25 GeV, but could be extended to 30 GeV at a later stage by adding RF equipment (missing-RF machine); to this end the magnet and vacuum system in the lattice are designed for 30 GeV operation. The insertion is foreseen for initial operation up to 25 GeV.

Colliding-beam e-p operation can take place during the SPS acceleration cycle over a proton energy range of 145 GeV to 400 GeV, and the proton beam can subsequently be used for fixed-target experimentation. Alternatively, intermediate plateaux in the SPS cycle could be used for e-p operation at fixed proton energy. The nominal proton energy of CHEEP has been chosen as 270 GeV, the maximum at which the present SPS can run in d.c. mode. The electron beam can be kept circulating for many hours.

At nominal proton energy the peak luminosity per intersection point is $0.5 \times 10^{32} \text{ cm}^{-2}\text{ per sec for an electron energy of 25 GeV. The corresponding centre-of-mass energy is 164 GeV, which compares well with the energy available in e-p options put forward by other laboratories.}
In the electron case, the crossing angle can be reduced during runs in which slow extraction is not required; this increases the luminosity to $0.75 \times 10^{32}$ cm$^{-2}$ sec$^{-1}$. The luminosity does not vary strongly with the proton energy. The electron energy can be varied over a wide range without appreciable loss in luminosity, e.g. at 13 GeV the latter is half of the nominal value.

Figure III.1 gives the basic layout of the two rings in the SPS tunnel. Initially, one interaction point is foreseen in the long straight section LSS 5 of the SPS which is entirely free of any special equipment. It may be possible to move the proton beam dump from the long straight section LSS 4 to LSS 1, thus making available another interaction point. A detailed study would be necessary in order to see how the proton beam dump could best be incorporated with the existing injection equipment in LSS 1.

An attractive feature of this scheme is that the high-intensity proton bunches only need to live over one, somewhat extended, synchrotron cycle and not over many hours, thus relaxing long-term stability requirements for intense, bunched proton beams. The electrons have a lifetime exceeding 10 hours, whereas radiative polarization of the electrons will build up with a time-constant of $\sim \frac{1}{2}$ hour. In order to take full advantage of polarized $e^\pm$, the insertion is designed in such a way that the polarization vector can be rotated into the longitudinal direction at the interaction point. Negative helicities are available over a wide range of electron (positron) energy, and positive helicity will be produced by spin-flip resonance crossing.

The protons are obtained by injecting three pulses from the CERN Proton Synchrotron (PS) into the SPS before acceleration. The electrons (positrons) require an injector synchrotron of 5 GeV. A possible location of the electron synchrotron on the CERN site is shown in Fig. III.2, which also illustrates the layout of the various transfer tunnels.

During the course of this study, many problems have been resolved, but a few have come to light for which we have not yet demonstrated satisfactory solutions. With the present parameters, the proton single-beam incoherent tune shift at an injection energy of 14 GeV in the SPS is rather large and a means of reducing it must be found. Possibilities exist in principle, but more study is required to determine their efficacy. The reduction in SPS flexibility resulting from the absence of slow extraction to the West Hall during $e^\pm$ physics is a regrettable operational constraint which we would like to alleviate without sacrificing the $e^\pm$ performance. Some parameters in the present design were of necessity frozen at an early stage; a subsequent revision could be expected to facilitate some of the engineering problems and to optimize the over-all performance of the facility.

2. **Performance**

The choice of 25 GeV for the nominal energy of the electron (positron) ring is dictated by the need to reach the highest possible c.m. energies for $e^\pm$ collisions within reasonable constraints of technology, finance, and operation. In particular, the RF system can be accommodated in one long straight section, and the RF power requirements are not excessive for a machine of this energy and luminosity. The magnet and vacuum systems can, for little extra cost, be designed for operation up to 30 GeV, leaving open the possibility of a future upgrading to higher energy by just adding an RF system.
The nominal proton energy is chosen as 270 GeV, the maximum at which the SPS can at present run in d.c. mode. With 25 GeV electrons the corresponding c.m. energy is 164 GeV. Higher c.m. energies up to 200 GeV, corresponding to 400 GeV protons in the SPS, can be obtained at a reduced duty cycle during pulsed SPS operation, either in flat-top or intermediate plateau mode.

The lowest proton energy at which synchronism between electron and proton bunches can be maintained is 145 GeV, owing to the protons not being fully relativistic. The range of proton energy is obtained from

\[ \gamma_p = \sqrt{\frac{R_p + \Delta R_p}{R_e - R_p}}, \]  

(III.1)

where \( R_p, R_e \) are the nominal orbit radii of protons and electrons, and \( \Delta R_p \) is the deviation from the nominal proton orbit. Assuming \( \Delta R_e \) to be limited to \( \pm 10 \) mm, synchronism between proton and electron bunches can be maintained over the proton energy range 145-400 GeV, with \( R_e \) exceeding \( R_p \) by 13 mm. The electron radius must be kept fixed to maintain horizontal damping.

The large energy ratio of protons and electrons makes it possible to use some common bending elements near the interaction region for beam separation, and still have the choice of positrons or electrons. This feature permits a rather small crossing angle for reaching high luminosity. The perturbation of the proton orbit due to field reversal of the common elements can be corrected locally, even at injection energy, within a reasonable aperture for the proton beam pipe. Since both electron and proton beams have smaller vertical than horizontal emittances, a horizontal crossing has been chosen.

The choice of parameters for maximum luminosity is a more complicated procedure for e-p than for e⁺-e⁻ machines because of the inherently different characteristics of the two beams. The luminosity \( L \) per interaction point is given by

\[ L = f_s \sum_{k_b} \frac{N_e}{k_b} \frac{N_p}{k_p} \frac{1}{F}, \]

where \( N \) is the total number of particles
\( k_b \) is the number of bunches
\( f_s \) is the revolution frequency.

The effective area \( F \) is given by

\[ F = 2\pi \left( \sigma_{zp}^2 + \sigma_{ze}^2 \right) \left( \sigma_{xp}^2 + \sigma_{xe}^2 + \left( \sigma_{sp}^2 + \sigma_{se}^2 \right) \phi^2 \right)^{\frac{1}{2}} \]

in the limit of

\[ \beta^* \gg \sigma^*_s, \quad \phi \ll 1, \]

with \( \sigma^* \) the r.m.s. beam size at crossing point
\( 2\phi \) the crossing angle
\( \beta^* \) the beta function at crossing point.

We have assumed the total number of protons to be \( N_p = 2 \times 10^{13} \) per pulse after the SPS improvement programme is implemented.
The number of electrons per bunch \( \frac{N_e}{k_b} \) is determined by the beam-beam tune shift\(^{11}\) tolerated by the protons, which has been assumed to be about 0.01, a factor of 2 above the generally accepted value for proton storage rings, since a proton lifetime of a few tens of seconds is sufficient to achieve a duty cycle approaching 100%.

The emittances of the proton beam are known, and those of the electrons are largely determined by the lattice and by an estimate of the optimum coupling. This together with the crossing angle defines the effective area \( F_e \), and the luminosity can then be calculated.

In order to restrict the number of electrons, which determines filling time and radiated power, it is desirable to operate with a minimum number of bunches \( k_b \). However, there is a lower limit to \( k_b \) set by the maximum tolerable number of protons per bunch \( \frac{N_p}{k_b} \), which is limited to about \( 3 \times 10^{11} \) by the RF system and by space-charge problems at injection. The beam-beam tune shift suffered by the electrons, usually the limit on \( \frac{N_p}{k_b} \), is below the admissible value; this is welcome as it helps to keep depolarization by beam-beam effects low. Finally, the number of bunches has to be adjusted to satisfy the numerology of a suitable PS-SPS transfer scheme.

In principle all the performance parameters are now determined, but some checks and adjustments have to be made to ensure that other constraints have been respected, such as performance of the injectors, stability criteria, the synchrotron radiation power loss of the electrons, etc. We then arrive at the performance figures forming part of the parameter list of Table III.1.

In the electron case, the crossing angle can be further decreased to \( 2\phi = 1.7 \) mrad during runs in which slow extraction is not required. For positrons, this is not possible with the present layout of the insertion. Figure III.3 shows the luminosity as a function of crossing angle, and it can be seen that the luminosity with the reduced crossing angle is increased to \( 0.75 \times 10^{32} \text{ cm}^{-2} \text{ sec}^{-1} \) at nominal energy.

It is proposed to collide the beam during an SPS cycle in order to minimize the interference with the fixed-target experiments. Excessive tune shifts in the proton beam will be avoided by separating the two beams vertically at low proton energies. Figure III.4 shows an example of an SPS cycle yielding a 25% duty cycle for e-p physics and providing the usual duty cycle for ejected beams. The dashed line indicates a cycle with a plateau at 270 GeV, giving a 40% duty cycle for e-p without substantially changing the duty cycle for fixed-target experimentation. These examples indicate that a fair margin is left to accommodate requirements which cannot be foreseen at present. The e-p duty cycle is essentially determined by the SPS fixed-target duty cycle and, therefore, by priority considerations.

In order to assess the merits of various modes of operation, it is important to know the variation of luminosity with proton energy when all other parameters are kept constant. The result is shown in Fig. III.5. The proton energy \( E_{p,\text{min}} \) defined as the energy where the beam separation is removed, appears as an important parameter. This is due to the fact that the beam-beam tune shift \( \Delta Q_p \) suffered by the proton beam is inversely proportional to the proton energy and proportional to the number of electrons. Since the tune shift should not exceed a given value \( (\Delta Q_p = 0.01) \), the number of electrons is determined by the lowest proton energy during the collision interval in the cycle; the intensity of the electron beam must be adjusted in such a way that the protons can just stand it at their lowest energy. Thus the electron intensity will not be optimum at higher energies in the cycle, which implies that the maximum luminosity cannot be reached at \( E_p > E_{p,\text{min}} \).
Table III.1

Basic parameters

<table>
<thead>
<tr>
<th></th>
<th>p-ring</th>
<th>e-ring a)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Luminosity (peak)</td>
<td>cm⁻² sec⁻¹</td>
<td>0.5 x 10¹²</td>
</tr>
<tr>
<td>Nominal energy</td>
<td>GeV</td>
<td>270</td>
</tr>
<tr>
<td>Total number of particles</td>
<td>2 x 10¹⁵</td>
<td>1.5 x 10¹³</td>
</tr>
<tr>
<td>Circulating current</td>
<td>mA</td>
<td>140</td>
</tr>
<tr>
<td>Number of bunches</td>
<td></td>
<td>60</td>
</tr>
<tr>
<td>Circumference</td>
<td>m</td>
<td>2π x 1100</td>
</tr>
<tr>
<td>Bending radius</td>
<td>m</td>
<td>741.3</td>
</tr>
<tr>
<td>Energy range for collisions</td>
<td>GeV</td>
<td>150-400</td>
</tr>
<tr>
<td>Injection momentum</td>
<td>GeV/c</td>
<td></td>
</tr>
<tr>
<td>Crossing angle</td>
<td>mrad</td>
<td>±2.5</td>
</tr>
<tr>
<td>Beta function βₓ/βₚ</td>
<td>m</td>
<td>6.5/0.6</td>
</tr>
<tr>
<td>Dispersion function Dₓ/Dₚ</td>
<td>m</td>
<td>0</td>
</tr>
<tr>
<td>Free space for detector</td>
<td>m</td>
<td>±5 m</td>
</tr>
<tr>
<td>Beam size at crossing αₓ</td>
<td>mm</td>
<td>0.34</td>
</tr>
<tr>
<td>Beam size at crossing αₚ</td>
<td>mm</td>
<td>0.072</td>
</tr>
<tr>
<td>Bunch length σₛ</td>
<td>mm</td>
<td>300</td>
</tr>
<tr>
<td>Energy spread cₑ/E</td>
<td></td>
<td>0.8 x 10⁻³</td>
</tr>
<tr>
<td>Beam-beam tune shift ΔQₓ/ΔQₚ</td>
<td></td>
<td>0.008/0.008</td>
</tr>
<tr>
<td>Lifetime</td>
<td>SPS cycle</td>
<td>20 h</td>
</tr>
<tr>
<td>Filling time (p e⁻/e⁺)</td>
<td>sec</td>
<td>1.8</td>
</tr>
<tr>
<td>Polarization time</td>
<td>min</td>
<td>-</td>
</tr>
<tr>
<td>Energy loss/turn</td>
<td>MeV</td>
<td>-</td>
</tr>
<tr>
<td>Total RF power</td>
<td>MW</td>
<td>51.8(96.7)</td>
</tr>
</tbody>
</table>

a) at 25 GeV operating with four orbit kinks

Using now Fig. III.5 to evaluate the luminosity averaged over the cycle shown in Fig. III.4 yields 0.75 x 10³¹ for the cycle indicated by the full line and 1.5 x 10³¹ for the cycle shown by the dashed line.

Examination of Fig. III.5 indicates that the luminosity, averaged over an SPS cycle, does not change drastically as function of Eₚ, min since the lower luminosity at low Eₚ, min is balanced by a longer collision interval in the cycle. If plateaux are used, the conclusions will obviously be different, and average luminosities close to the peak luminosity of the plateau energy can be achieved.

In order to explore various kinematical regions, the electron energy would be varied as well. Figure III.6 shows the luminosity plotted versus electron energy for the case where kinks are used to keep the polarization time at 30 min at all electron energies. The proton energy is 270 GeV in all cases.
It is instructive to investigate how the performance would be affected by a change in certain parameters which cannot be predicted with certainty. Figure III.7 shows the influence of the proton bunch length, which is assumed to be blown up by a longitudinal instability from its theoretical minimum of 16 cm to 30 cm. An improvement would be possible if this blow-up could be limited. Another parameter which is hard to assess is the coupling in the electron ring; it determines the ratio of vertical to horizontal emittance. It can be concluded from Fig. III.8 that the coupling is not a critical parameter in this machine.

The lifetime of the proton beam is hardly significant given the length of the cycle. On the other hand, proton losses in the insertion must be very small to reduce background, and an ultra-high vacuum comparable to the ISR is therefore required in this region. Different considerations apply for the electron ring where the beam should have a long lifetime, particularly because of the relatively long filling time of the positrons. The average pressure in the electron ring is around 10⁻⁹ Torr, providing a bremsstrahlung lifetime of 30 hours. The overvoltage of the RF has been chosen to be such that the lifetime due to quantum fluctuations approaches 100 hours. The losses brought about by beam-beam interactions are negligible; the strongest effect is beam-beam bremsstrahlung, giving a lifetime of 290 hours. Thus a lifetime of about 20 hours can be expected, nearly independent of energy as explained in Section 5.4. This is convenient not only for colliding beam operation but also for injection and acceleration.

The degree of longitudinal polarization at the crossing point is greater than 76% over the range 15 to 25 GeV. The layout of the interaction region is such as to produce a negative natural helicity; this is the preferred sign for many experiments. The opposite helicity will be produced by spin-flip; the efficiency of this process is not yet well quantified and needs further investigation.

The constraints imposed by fixed-target operation of the SPS and the desire to keep modifications in the SPS to a minimum have adverse effects on the performance. Since it is interesting to find out by how much the performance is affected by these constraints, we try to assess roughly what improvement would be possible if the e-p programme had such high priority that these restrictions would be removed.

For this purpose we assume that a proton RF system for very high intensity operation, as proposed for the DESY facility 12), would be available, which could handle \( 8 \times 10^{11} \) protons per bunch. At these high intensities, proton injection at 14 GeV/c could no longer be maintained, but the injection energy would have to be raised to 26 GeV, the highest PS energy, in order to keep space-charge effects under control and to avoid crossing transition in the SPS. The number of electrons could be increased by 2%, raising \( \Delta q_p \) from 0.008 to 0.01; the required RF power would augment by 1.5 MW to 9.5 MW. Then the electron-proton luminosity would reach \( 2.5 \times 10^{32} \) cm⁻² sec⁻¹, and \( 1.5 \times 10^{32} \) cm⁻² sec⁻¹ could be obtained for positron-proton collisions, both values quoted at nominal energy.

In order to have maximum luminosity also at lower centre-of-mass energies, it is proposed to lower the electron energy and not the proton energy, keeping the number of bunches \( k_b \) at the maximum given by the available RF power. Obviously, the number of bunches cannot be varied continuously, but in steps satisfying the numerology set by the proton injector. Using 140 bunches at 20 GeV yields a luminosity 25% higher than the luminosity at 25 GeV. This increase becomes 40% at 18 GeV where 220 bunches would be used.
For a dedicated facility, more than two interaction regions might be offered, and one would certainly try to boost the energy at which the SPS can run in d.c. mode closer to 400 GeV.

However, all this has not been studied in any detail since these considerations are beyond the scope of this report, which is based on the fact that a vigorous fixed-target programme will exist for the SPS in the next decade.

3. INTERACTION REGION GEOMETRY AND OPTICS

3.1 Crossing region

CHEEP differs fundamentally from earlier proposals for e-p studies in that the SPS will continue to function as an accelerator rather than a storage ring. Thus the insertion must accommodate the cycling of the proton energy without disturbing the stored electron (or positron) beam. For high luminosity a small (or zero) crossing angle is desirable, but this makes separation of the two beams difficult. One approach is to use a very strong common magnet which bends the lower energy beam away from the higher energy beam. Separate bending magnets are then used for the two beams to correct the effects of the large orbit distortion. The beam separation must therefore provide sufficient space for the magnet yokes, and the high fields required produce an intense synchrotron radiation background in the detectors which would make experimentation difficult if not impossible. Another approach is to use septum magnets to bend one beam away from the other. The beam sizes at the septum and the septum thickness then determine the minimum crossing angle and hence the luminosity. For CHEEP a new configuration is proposed using a combination of these two methods. There are two common fixed-field bending magnets close to the crossing point. As the proton energy varies during a machine cycle from 14 GeV to 400 GeV, the deviation of the proton beam also varies during the cycle. Since it is important that the protons pass centrally through the strong low-beta quadrupoles, the orbit distortion of the proton beam due to the common bending magnets must be corrected before reaching the quadrupoles. This can be achieved using two double-sided septum magnets having different fields on either side of the septum. The electrons (or positrons) see a fixed field equal to that in the common bending magnets. For a given beam-beam separation required at the quadrupoles, this configuration gives the smallest possible amount of synchrotron radiation. Even so, great care must be taken in the design of the collimators and the vacuum chamber in this region (see Section 2 of Chapter II). The protons also see a constant field in the septum magnets (i.e. the field does not vary with proton energy and the septum magnets are not of the pulsed type), and these fields are so arranged that they exactly compensate, in position and angle, the effect of the common bending magnets. The sign of the field in the common magnets depends on whether electrons or positrons are used in the CHEEP ring, and so two different field settings must be employed with different proton trajectories, as shown schematically in Fig. III.9. In the electron case an additional crossing angle is introduced. Table III.2 gives the parameters of the separating magnets, while Figs. III.10 and III.11 show the trajectories for 25 GeV electrons or positrons and 14 GeV protons. The centre lines are indicated by a dot-dash and show that the protons are displaced a maximum of 36.4 mm in the positron case and 30.6 mm in the electron case. This extra aperture requirement is the price to be paid for the additional separation obtained. Also shown in Figs. III.10 and III.11 are the beam envelopes (dashed lines) calculated for ±10 σ for the electrons and ±6 σ for the protons. Finally the closed-orbit error is added (solid lines) scaled from the measured SPS value (5.0 mm peak-to-peak). The
rather large separation of the beams at the position of the septa is clearly seen, but this is not the limiting case. Slow ejection of the protons at high energy also requires a large aperture and in this case the proton orbit bump is much smaller. This problem is dealt with in Section 6 of this chapter. The main advantage of this configuration is that the innermost septum has been placed well back where the electrons (or positrons) have been bent away from the protons by the common bending magnets.

Under certain conditions it is possible to reduce the crossing angle below 5 mrad. It is a basic design feature of the crossing region that the septum magnets are movable, so that the septa can be retracted when the SPS is not used for e-p collisions. Thus it will be possible to vary the operating positions of the septa and the fields in the magnets independently, which offers a range of configurations in those runs where slow extraction from the SPS is not required. There are some constraints, however. Both beams must still pass centrally through their respective low-beta quadrupoles, and this fixes the fields on both sides of the septum for a given crossing angle and common field. Also, any reduction in the crossing angle naturally requires an increased common field to give the necessary separation at the first septum. In the positron case, the correcting field for the protons is already almost at the limit in the first septum magnet (see Section 3.5), so very little improvement is possible. In the electron case, the main limitation is the synchrotron radiation background. One example of many is shown in Fig. III.12, where a crossing angle of 1.7 mrad is obtained giving a luminosity of $0.75 \times 10^{32} \text{ cm}^{-2} \text{ sec}^{-1}$. The maximum field is exactly twice that in the standard configuration, doubling the critical energy of the synchrotron radiation. While the background problems are increased, they do not appear to be prohibitive. However, about 20% of the radiation strikes the downstream septa, which will require additional cooling.

**Table III.2**

Parameters of separating magnets

<table>
<thead>
<tr>
<th>Element</th>
<th>Length</th>
<th>Field</th>
<th>Angle (mrad)</th>
<th>Field</th>
<th>Angle (mrad)</th>
<th>Field</th>
<th>Angle (mrad)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>m</td>
<td>T</td>
<td>25 GeV</td>
<td>T</td>
<td>14 (270) GeV</td>
<td>T</td>
<td>14 (270) GeV</td>
</tr>
<tr>
<td>Drift space</td>
<td>5.5</td>
<td>-</td>
<td>5.0 (crossing)</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-4.931 (-0.256)</td>
</tr>
<tr>
<td>Common magnet</td>
<td>1.9</td>
<td>0.165</td>
<td>3.75</td>
<td>-0.165</td>
<td>-6.696 (-0.347)</td>
<td>0.165</td>
<td>6.696 (0.347)</td>
</tr>
<tr>
<td>Drift space</td>
<td>0.3</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>Common magnet</td>
<td>1.9</td>
<td>0.165</td>
<td>3.75</td>
<td>-0.165</td>
<td>-6.696 (-0.347)</td>
<td>0.165</td>
<td>6.696 (0.347)</td>
</tr>
<tr>
<td>Drift space</td>
<td>0.3</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>First septum</td>
<td>1.9</td>
<td>0.165</td>
<td>3.75</td>
<td>0.823</td>
<td>33.482 (1.756)</td>
<td>-0.104</td>
<td>-4.231 (-0.219)</td>
</tr>
<tr>
<td>Drift space</td>
<td>0.3</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>Second septum</td>
<td>1.9</td>
<td>0.165</td>
<td>3.75</td>
<td>-0.494</td>
<td>-20.089 (-1.042)</td>
<td>-0.104</td>
<td>-4.231 (-0.219)</td>
</tr>
</tbody>
</table>
3.2 Polarization rotation

There is great interest in producing longitudinal polarization of the electrons (or positrons) in e-p collisions\(^1\)\(^6\), and this can be obtained by rotating the naturally produced transverse polarization (see Section 4 of this chapter) into the longitudinal direction by a series of vertical and horizontal bends\(^1\)\(^1\)\(^7\). A detailed analysis of the scheme proposed for CHeEP\(^1\)\(^8\) shows that for fixed geometry the transformation of the polarization vector is strongly energy dependent and can only be optimized at one energy. A scheme in which the bending angles in the rotation magnets are modified depending on the working energy is therefore proposed\(^1\)\(^5\).

The rotation of the polarization vector is obtained by a sequence of vertical (RV), horizontal (RH), vertical (V), and horizontal (H) bends (Fig. III.13). The two vertical bends are equal and opposite with bending angle \(\alpha\) and precession angle \(\psi\) given by

\[
\psi = \gamma \left( \frac{\frac{g-2}{2}}{g} (\frac{g-2}{2}) \right) = 2.2734 \text{E(GeV)} \alpha,
\]

where \(g\) is the gyromagnetic ratio. The reverse horizontal bend (RH) with bending angle \(\beta\) and precession angle \(\theta_1\) is greater than the horizontal bend (H) by an amount equal to the crossing angle \(2\phi\) (5 mrad for CHeEP), and the precession angle \(\theta_2\) is given by

\[
\theta_2 = -\left( \theta_1 - \gamma \left( \frac{\frac{g-2}{2}}{2} \right) \right) 2\phi.
\]

The longitudinal projection of the initially transverse polarization vector is given by\(^1\)\(^8\)

\[
P = -\sin \theta_2 \sin \theta_1 \sin \psi + \cos \theta_2 (1 - \cos \theta_1) \sin \psi \cos \psi.
\]

It is possible to solve this equation at any energy to obtain values of \(P = +1\) or \(-1\) corresponding to the "normal" and "reverse" polarization directions, with the result shown in Fig. III.14. The range of bending angle required is considerable, and geometrical limitations prevent the "reverse" polarization case from being a viable solution for CHeEP.

The variation of bending angle is obtained by splitting up each of the rotation bends into three magnets. In the standard configuration all three are equally powered. By varying these fields independently it is possible to alter the total bending angle required to give the same transverse displacement of the beams. The innermost horizontal bend H is largely made up of the common and septum magnets described above, and the geometry in this region must remain constant. This is mainly to ensure that the electron beam passes centrally through the low-beta quadrupoles but also to facilitate the design of the synchrotron radiation absorbers required to shield the intersection region. Thus only a small fraction of H can be varied. The details of the optimization of the bending angles is rather tedious and will not be carried out here (see Ref. 15 for details). For CHeEP, a standard geometry has been adopted with all rotation magnets producing 10 mrad bending. The other extreme is at low energy (15 GeV) when the outer magnets of each set are deactivated and the inner magnets produce about 42 mrad each. Thus the critical energy of the synchrotron radiation remains approximately constant (124 keV at 25 GeV, 112 keV at 15 GeV). The layout of half the insertion is shown in Fig. III.15; the other half is an antisymmetric image. The additional
aperture requirement in this region is 130 mm horizontally and 80 mm vertically. The resulting projection of the transverse polarization in the longitudinal direction is over 94% over the whole energy range of interest, 15-25 GeV.

3.3 Transition regions

As can be seen from Fig. III.15, the polarization rotation bends produce a vertical and horizontal displacement of the electron beam relative to the proton beam, which is straight in this region. In order to bring the electrons to their standard positions 1.1 m directly above the proton ring, two transition regions are required in the arcs immediately adjacent to the long straight section. On the entrance side (for the electrons) the electron ring is 0.645 m above, and 1.024 m outside the proton beam. The horizontal separation is obtained by introducing a straight section into the electron lattice and making up the required bending angle with standard lattice magnets run at twice the field strength. A combination of horizontal and vertical bends produces spin precession as shown above. Since the spin is in the correct orientation at either end of the central region, the vertical bends required to bring the beams to the correct height should not disturb this. A sufficient condition is that the vertical bends be placed in the straight section inserted into the electron lattice. On the exit side the electron ring is 0.645 m below and 1.024 m inside the proton beam. A strong horizontal bend is introduced into the electron ring immediately following the rotation bends to separate the electron and proton rings in the arc. The electron ring goes straight along a chord before rejoining the proton ring. This very long straight section is packed full of vertical bending magnets required to make up 1.745 m vertical displacement. The layout of the whole insertion is shown later in Fig. III.50.

It is worth noting that while the over-all design is fixed by the tunnel constraints, the detailed layout is still subject to minor changes. When exact details of all the magnets are known, some slight alterations may in fact be necessary.

3.4 Betatron matching

The characteristics of the interacting region determine the physics potential of the whole facility. In order to obtain the desired luminosity, low values of the vertical and horizontal beta-functions are required at the interaction point for both electrons and protons. Since the emittance of both beams is smallest in the vertical direction, a horizontal crossing angle was adopted. The luminosity is then strongly dependent on the vertical beta-function, which should be as low as possible.

The proton ring is the easier of the two. All elements are (nominally) in the horizontal plane, making the vertical dispersion zero to first order. The horizontal dispersion enters only weakly into the luminosity formulae and, while it should be small, it does not necessarily need to be zero. A vertical beta-value of 0.6 m and a horizontal beta-value of 6.5 m with zero horizontal dispersion have been obtained (Fig. III.16). For this, four special quadrupoles either side of the crossing point are used, and small adjustments are made to four quadrupoles in the lattice. It is possible that the addition of small auxiliary tuning quadrupoles at these places (rather than auxiliary windings on modified SPS quadrupoles) would be more suitable for practical reasons, and future work will concentrate on this possibility. The insertion may be retuned to obtain a reasonable facsimile of the SPS normal structure (Fig. III.17) which, coupled with a retraction of the septa, could return the SPS almost to its present state when not used for e-p physics.
The electron ring is considerably more complicated, and at present no complete solution, meeting all of the requirements simultaneously, has been obtained. The vertical bending in the interaction region adds (literally!) another dimension to the matching problem. Vertical dispersion in the lattice produces depolarization (Section 4.2.5 of this chapter), which is undesirable. Vertical dispersion in the intersection region enlarges the beam height, which reduces the luminosity. In addition, the geometry in the transition regions is different on the two sides of the insertion so two separate matches must be obtained. The matching obtained for the entrance side is shown in Fig. III.18. It is fully matched to the lattice on the right-hand side. Zero vertical dispersion is obtained at both ends and zero horizontal dispersion at the intersection. A vertical beta-value of 0.3 m and a horizontal beta-value of 1.5 m have been obtained. The high vertical beta-values in the polarization rotation bends are undesirable and efforts are being made to lower them. For the exit side the best match so far obtained is shown in Fig. III.19. This is completely matched both into the lattice and into the other half of the insertion, except for a small (5 cm) vertical dispersion which has not yet been removed. Some minor changes to the geometrical layout in this region should solve this difficulty.

Chromaticity correction of the two insertions has not yet been attempted and, owing to lack of man-power, will probably not be started until a later stage of the study.

3.5 Special components

The layout of the interaction region involves, of necessity, many non-standard magnets and these are listed in Table III.3. Although many of these are fairly ordinary there are some which must satisfy stringent requirements, and these magnets are discussed in detail below. The magnetic field characteristics were calculated using the FATIMA program\textsuperscript{13} assuming that the magnetic properties of the iron are similar to those of the SPS magnets.

3.5.1 Common bending magnets

In principle these magnets are simple, low-field magnets but they require considerable modification to meet the special requirements of the experimenters. Downstream of the insertion the magnets have a normal yoke with a large aperture for the passage of the synchrotron radiation (Fig. III.20), as described in more detail in Section 2 of Chapter II. Upstream of the intersection the common magnets (and also the double septum magnets) play a vital second role as integrated parts of a hadronic shower counter for all physics experiments. Thus these magnets have a smaller aperture, and the yokes consist of alternate layers of iron and scintillator approximately 3 mm thick and perpendicular to the beam direction. Then assuming coils with a 70% filling factor, the good field region ($\Delta B/B < 10^{-7}$) extends to within 15 mm of the iron or coils. The field variation near the iron caused by the laminated construction is periodic in the longitudinal direction and has little influence on the field integral. To improve the detection efficiency the coils should be made as thin and transparent to particles as possible.

3.5.2 Double-sided septum magnets

The double-sided septum magnets are similar to the common magnets. Again two different types are required: one with a compact yoke and a "wide" aperture, and one with a laminated yoke, for the insertion of scintillators as described in Section 3.5.1, and a small aperture.
However, three coils and two opposed field levels are required for the bending of the electron beam and for the correction of the proton orbit. The arrangement of the coils and the flux lines are shown in Fig. III.21. In the case presented, the central conductor carries a high current density of 60 A/mm² on the average, which is, however, still technically possible. The thickness of the septum is determined by the beam separation at the entrance of the magnet and can be increased towards the other side by ~ 50% if necessary. The forces on the septum can be expressed by a pressure in the direction of the lower field side of $1.7 \times 10^5 \text{ N/m}^2$, which can be taken by a suitably designed vacuum chamber on the weak-field side. The magnetic properties are similar to those of the common magnets since the iron is not saturated even in the laminated yoke at maximum field. (Maximum induction is 1.7 T in the iron.)

### 3.5.3 First electron quadrupole $Q_1$

Only 120 mm, of which 45 mm have to be useful aperture, separate horizontally the central orbit of the electron beam from the edge of the proton beam. The most suitable type of quadrupole in this case seems to be the modified "Panofsky" quadrupole described by Morpurgo. In this type of lens the useful elliptical aperture is delimited by iron, shaped to follow an equipotential curve (e.g. AB in Fig. III.22) and by the conductors. Thus the field quality is determined by the shape of the iron (in the following assumed to be perfect) and the homogeneity and shape of the coils. Conductors shaped as shown in Fig. III.22 produce a gradient error of less than $2 \times 10^{-4}$ at 15 mm and of $6 \times 10^{-5}$ at 10 mm away from the coil (the precision of the computation is better than $10^{-4}$). The current density is 25 A/mm² on the average. With an aperture of ±60 mm horizontally and ±80 mm vertically this quadrupole still fits into the given space and produces the necessary gradient with sufficient quality.

### 3.5.4 Proton quadrupoles

For the first and second proton quadrupoles there is only slightly more space available horizontally than for the first electron quadrupole. However, a higher gradient is required. As the type of quadrupole described above in Section 3.5.3 is rather insensitive to saturation effects, this type is also proposed for the first two proton quadrupoles. The next two quadrupoles downstream of the interaction area are more conventional.

### 3.5.6 Other special features

An ultra-high vacuum is required in the whole insertion region to limit the background from beam-gas events. This will be achieved entirely with lumped pumps since all the elements are short and have a rather large aperture. The vacuum chamber at the intersection itself must be designed in conjunction with the detector. Provisions should be made for baking this part of the vacuum system up to 300°C. The problem of higher mode losses in such vacuum chambers has not been studied in detail, but it is possible that additional cooling may well be required in some regions. The collimators for the absorption of the synchrotron radiation are fitted into the gaps between the upstream magnets, and are described in detail in Section 2 of Chapter II.
Table III.3
Special magnets in the insertion

<table>
<thead>
<tr>
<th>Magnet function</th>
<th>Max. field T</th>
<th>Length m</th>
<th>Aperture (mm)</th>
<th>Free space between the beam mm</th>
</tr>
</thead>
<tbody>
<tr>
<td>Common bending magnet</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>upstream</td>
<td>0.165</td>
<td>1.9</td>
<td>130</td>
<td>450</td>
</tr>
<tr>
<td>downstream</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Double-sided septum</td>
<td>0.823</td>
<td>1.9</td>
<td>130</td>
<td>490</td>
</tr>
<tr>
<td>upstream</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>downstream</td>
<td></td>
<td></td>
<td></td>
<td>20</td>
</tr>
<tr>
<td>Low-beta quadrupoles</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>electron beam: Q₁</td>
<td>10.1 T/m</td>
<td>1.0</td>
<td>130</td>
<td>90</td>
</tr>
<tr>
<td>Q₂</td>
<td></td>
<td></td>
<td></td>
<td>75</td>
</tr>
<tr>
<td>proton beam: Q₁</td>
<td>16 T/m</td>
<td>5.5</td>
<td>120</td>
<td>70</td>
</tr>
<tr>
<td>polarization rotation</td>
<td></td>
<td></td>
<td></td>
<td>100</td>
</tr>
<tr>
<td>horiz.</td>
<td>0.3</td>
<td>2.8</td>
<td>190</td>
<td>250</td>
</tr>
<tr>
<td>vert.</td>
<td>0.3</td>
<td>2.8</td>
<td>190</td>
<td>250</td>
</tr>
<tr>
<td>Orbit restoring</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>horiz.</td>
<td>0.22</td>
<td>6.26</td>
<td>190</td>
<td>250</td>
</tr>
<tr>
<td>vert.</td>
<td>0.36</td>
<td>2.8</td>
<td>190</td>
<td>250</td>
</tr>
</tbody>
</table>

4. POLARIZATION OF THE ELECTRON (POSITRON) BEAM

An important feature of an e-p facility is to have polarized electron (positron) beams; this has been a basic assumption of Llewellyn Smith and Wiik\textsuperscript{16}). The design concepts of CHEEP therefore put a strong emphasis on obtaining a high degree of polarization and an acceptably short polarization time over as wide an energy range as possible. Furthermore, the layout of the interaction regions (Section 3 of this chapter) is specifically chosen to rotate the polarization vector into the longitudinal direction at the crossing points\textsuperscript{18}) using a scheme analogous to that proposed earlier for the LSR e-p option. Methods are also foreseen for providing either positive or negative helicities for the interactions.

In the range of CHEEP energies there are several potential sources of depolarization which must be very carefully controlled. To monitor the effects of the various corrections will require a rapid on-line facility for measuring the degree of polarization and its variation with time. Special provision is therefore made to incorporate a sensitive, high-resolution polarimeter using back-scattering of a circularly polarized laser beam.

4.1 Radiative polarization rate

For a storage ring with an isomagnetic field the polarization time $\tau_p$ is given by\textsuperscript{21})

$$\frac{1}{\tau_p} = \frac{5\sqrt{2}}{8} \frac{\mu R}{m_0 R^3} \gamma^5,$$  \hspace{1cm} (III.2)
where R is the average radius and $\rho$ the bending radius of the ring; $m_0$ and $r_0$ are the electron rest mass and classical radius, respectively; $\hbar$ is the reduced Planck's constant; and $\gamma$ is the Lorentz factor. For CHEEP at 25 GeV the polarization time would be about 100 minutes and would increase rapidly at lower energies if no special measures were taken.

In the presence of depolarizing effects, the polarization time also determines the limiting degree of polarization achievable. Thus in order to obtain highly polarized beams in a sufficiently short time and over a wide range of operating energy, CHEEP will be equipped with a system of short, high-field magnets ("kink magnets") designed to maintain an approximately constant polarization time of $\sim 50$ min over an energy range of 15-25 GeV. Compared with a more conventional wiggler-magnet system, the proposed arrangement, described in Section 5 of this chapter, has the advantage of not reducing the asymptotic degree of polarization, since no reverse bending is used.

4.2 Depolarization effects

Transverse radiative polarization in electron (positron) storage rings was predicted theoretically by Sokolov and Ternov and has been experimentally observed at ACO, VEPP-2, and in the early operation of SPEAR I. This natural polarization process can be disturbed by various depolarizing effects, most of which are associated with spin-precession resonances of the form

$$\gamma_a = k + k_z Q_z + k_x Q_x + k_s Q_s.$$  \hspace{1cm} (III.3)

Here $\gamma$ is the Lorentz factor; $a = (g-2)/2$ is the anomalous part of the gyromagnetic ratio; $k, k_z, k_x, k_s$ are integers (positive or negative); and $Q_z, Q_x, Q_s$ are the wave numbers for vertical, horizontal, and energy oscillations, respectively.

The most important spin resonances in high-energy machines are those identified by $k$ ("imperfection" resonances) and $k_z$ ("intrinsic" resonances) in Eq. (III.3), and driven by harmonics of radial components of the magnetic fields. Longitudinal field components are much less important in CHEEP. The radial fields can arise, for example, from vertical motion in quadrupoles due to closed orbit errors or vertical betatron oscillations, and lead to perturbed motion of the spin precession. If the operating energy is chosen so that $\gamma_a$ is not too close to a resonance, the precession perturbations can usually be kept rather small.

In electron machines at high energies, however, a quantum emission causes an abrupt change in trajectory parameters without change of spin direction at this instant. The subsequent trajectory oscillation gives rise to a precession perturbation but, whereas the initial oscillation disappears owing to radiation damping, the perturbed precession remains, with a consequent loss of correlation between trajectory and spin motions. In an ensemble of many particles this process results in stochastic depolarization, also called spin diffusion, and can extend the influence of resonances outside their immediate vicinities.

The vertical motion with its associated radial fields can come about from several sources such as direct recoil from quantum emission, vertical closed-orbit errors, vertical betatron oscillations excited by coupling from the horizontal plane, vertical dispersion, beam-beam forces at the interaction region, etc. In general these sources have different depolarization characteristics and it is necessary to consider them separately.
A measure of such effects is the depolarization time constant $\tau_d$ associated with one or more of the mechanisms involved. With a finite value of $\tau_d$ the evolution of the beam polarization $P$ follows the equation

$$\frac{dP}{dt} = -\left(\frac{1}{\tau_p} + \frac{1}{\tau_d}\right) P + \frac{8}{5\sqrt{3}} \frac{1}{\tau_p},$$

with solution

$$P = \frac{8}{5\sqrt{3}} \frac{\tau_d}{\tau_d + \tau_p} \left\{1 - \exp\left[-\left(\frac{1}{\tau_p} + \frac{1}{\tau_d}\right) t\right]\right\},$$

In the absence of depolarizing effects the limit $P(t \to \infty) \approx 92.4\%$; otherwise it is reduced by a factor $\tau_d/(\tau_d + \tau_p)$.

In subsequent sections we make approximate estimates of the most important depolarizing sources in CHEEP using simplified formulae which are valid strictly only for "smooth" machines. More reliable and precise calculations, taking into account the detailed lattice properties, require more sophisticated analyses and computer programs, which are not yet available. It is normally sufficient to evaluate the depolarization in CHEEP at the nominal energy of 25 GeV, since the effects become generally weaker at lower energies.

4.2.1 Direct vertical quantum excitation

The recoil momentum on an electron due to the emission of a synchrotron-radiation photon has in general a vertical component which gives rise to the "natural" beam size vertically. Since the photon emission takes place in a very short time interval, of the order of $10^{-11}$ sec, the resultant perturbation of the orbit is abrupt. In terms of Fourier components the orbit perturbation has a very broad spectrum which can excite a large number of resonances. On the other hand, because the recoil momentum from the photon is small compared with the electron momentum, the strength of the resonance excitation is also small.

An approximate formula for a smooth machine can be derived from the results of Derbenev and Kondratenko:

$$\frac{1}{\tau_d} \approx \frac{(\gamma a)^2 Q_z^2}{(\gamma a + k + Q_z^2)^2} \frac{\pi}{R} \frac{1}{2\tau_E},$$

where $\lambda$ is the Compton wavelength, $R$ is the machine radius, and $\tau_E$ is the energy damping time. If we suppose that $(\gamma a)$ is only one-tenth of an integer away from the resonance condition, then with $Q_z \approx 28.2, \gamma a \approx 57$ at 25 GeV, $R = 1100$ m, $\tau_E \approx 12 \text{ msec}$, and $\lambda = 3.8616 \times 10^{-13}$ m, the depolarization time $\tau_d$ is 73 h from this source. This is a rather safe figure since, although a form factor would have to be applied to Eq. (III.6) to correct for a "non-smooth" machine, we have assumed an unnecessarily small value for the resonance denominator.

In fact, from the form of this denominator, it is evident that the resonances occur in pairs, separated by twice the non-integral part of $Q_z$. Thus, in order to make available as much free space as possible between adjacent pairs it is advantageous to operate with a small non-integral part of $Q_z$. 
In a perfect machine with superperiodicity S, only those values of k which are multiples of S would be significant in Eq. (III.6). Since GHEEP will have S = 1 we must suppose that all integers k are equally important.

4.2.2 Depolarization from vertical closed-orbit distortions

The depolarization rate arising from the kth Fourier harmonic \( z_k \) of the vertical closed-orbit (c.o.) distortion is given in Ref. 26 as

\[
\frac{1}{\tau_d} = \frac{1}{2} \frac{1}{\gamma^2} \frac{d}{dt} \langle \delta \gamma^2 \rangle \left( \frac{Q_x^h}{(\gamma a - k)^h} \right) \left( \frac{Q_x^h}{((\gamma a - k)^2 - Q_x^h)^2} \right) \frac{z_k^2}{R^2},
\]  

(III.7)

where R is the machine radius. Here,

\[
\frac{1}{\gamma^2} \frac{d}{dt} \langle \delta \gamma^2 \rangle = \frac{55}{24 \sqrt{5}} \frac{r_0 c \pi \gamma^5}{R \rho^2}
\]

(III.8)

is the quantum excitation rate, with \( r_0 \) the classical radius and \( \pi \) the Compton wavelength of the electron, and with \( \rho \) the bending radius.

The behaviour of formula (III.7) is dominated by the resonance denominator \((\gamma a - k)^h\). If we assume a smooth Fourier spectrum of the closed-orbit harmonics, resonances occur for all integers k and are separated in energy by \( m_0 c^2/a = 440 \text{ MeV} \). To minimize the depolarization, the energy is chosen to make \( \gamma a \) half integral and then \((\gamma a - k)^-h = 16\). The factor containing \( Q_x \) takes account of quantum jumps in horizontal betatron motion; it is essentially unity for \( Q_x >> 1 \) as in GHEEP.

The harmonic amplitude \( z_k \) may be related approximately to the peak c.o. amplitude \( \zeta \) by considering a regular machine with normal error statistics. Following Guignard\(^{27}\) we write

\[
z(\phi) = \sum_{k=1}^{\infty} \frac{Q_z^2}{Q_z^2 - k^2} \left[ f_k \cos k\phi + g_k \sin k\phi \right].
\]

(III.9)

For a machine with \( M \) magnet elements the spectral density of \( f_k, g_k \) is substantially uniform up to \( k = M \) and we can write \( f_k = f, g_k = g \). The summation over the \( \cos k\phi \) terms can then be expressed exactly by

\[
Q_z^2 f \sum_{k=1}^{\infty} \frac{\cos k\phi}{Q_z^2 - k^2} = \frac{\pi}{2} \frac{Q_z^2 \cos (Q_z (\pi - \phi))}{\sin \pi Q_z} - 1,
\]

(III.10)

in which the first term clearly dominates for \( Q_z >> 1 \). There appears to be no closed expression for the \( \sin k\phi \) series so we invoke the statistical nature of the errors by putting \( g = f \), assuming the same contribution from \( \sin k\phi \) and \( \cos k\phi \) terms and adding these in quadrature. Putting \( \cos (Q_z (\pi - \phi)) = 1 \) then gives

\[
\zeta = \frac{f \pi Q_z}{\sqrt{2} \sin \pi Q_z},
\]

(III.11)

From Eq. (III.9) the \( k \)th harmonic amplitude is, similarly,
\[ z_k = \frac{\sqrt{2} f Q^2_x}{Q^2 - k^2}, \]

and hence
\[ z_k = 2 \frac{Q^2_x \sin \pi Q^2_x}{\pi (Q^2_x - k^2)}. \]  

(III.12)

Replacing Eq. (III.12) in (III.7) and omitting the \( Q_x \) factor \((\sim 1)\) gives
\[ \frac{1}{\tau_d} = \frac{2}{\pi^2} \frac{1}{\gamma^2} \frac{d}{dt} \left( \delta \gamma^2 \right) \frac{(\gamma a)^{n} k^n}{(\gamma a - k)^n} \frac{Q_x^2 \sin^2 \pi Q^2_x}{R^2} \frac{Q^2_x}{(Q^2_x - k^2)^2}. \]  

(III.13)

For CHEEP at 25 GeV the quantum excitation rate from Eq. (III.8) is about \( 2 \times 10^{-6} \) sec\(^{-1}\) and \( \gamma a = k = 57 \) for the nearest resonance. Then, assuming \( Q^2_x = 28.2 \), the depolarizing time corresponding to a peak vertical closed-orbit amplitude \( z = 2.5 \) mm would be \( \tau_d \sim 16.3 \) hours, which is certainly adequate. However, although a peak c.o. amplitude of 2.5 mm or so is currently obtained in the SPS, the closed-orbit control in CHEEP, necessary to reach the highest polarization levels, may have to be somewhat tighter than the above figures indicate. This is because the approximate formulae are probably somewhat optimistic when applied to CHERP; in particular additional resonance terms can arise from the periodicity of the alternating gradient focusing.

4.2.3 Vertical excitation from coupling

In practical storage rings with alignment errors, a large fraction of the vertical betatron amplitude is produced by coupling from motion in the horizontal plane. Although the vertical beam size due to coupling is normally very much greater than the "natural" size determined by direct vertical quantum excitation, the excitation of spin resonances by coupling is intrinsically much weaker.

The essential qualitative feature of this process is the rate of energy transfer from horizontal to vertical motion which, for a few tenths of an integer tune split, takes place in a few tens of revolutions. This is slow compared with the precession rate \( \gamma a \) (several tens per revolution) but fast compared with the damping rates \((\sim 10^{-3} \) per revolution\). Thus the precession perturbations build up and decay adiabatically, and are in close correlation with the trajectory parameters at all times, provided \( \gamma a \) is not very close to a spin resonance.

The contribution to the depolarization rate from coupled betatron motion is given, in smooth approximation, by Derbenev (private communication) as
\[ \frac{1}{\tau_d} = \frac{(z^2)}{R^2} \frac{16(\gamma a)^2 \sin^2 \left( \pi (Q^2_x - Q^2_z) \right)}{Q_x (\tau_z + \tau_X)}, \]  

(III.14)

where \( (z^2) \) is the square of the vertical beam size due to coupling, and \( \tau_z, \tau_X \) are the vertical and horizontal damping times, respectively. The expression is valid not too close to a resonance of type \( \gamma a + k \pm Q_X = 0 \). It would appear from Eq. (III.14) that this source of depolarization vanishes for \( Q^2_x - Q^2_z = \) integer; this is not strictly true but results from the nature of the approximation.
Assuming for CHEEP a fractional Q-split of 0.2, $\tau_\chi + \tau_\kappa \approx 48$ msec, $Q_e \approx 28$, $\gamma_a \approx 57$, and $(z^2) \approx (4 \text{ mm})^2$ corresponding to $\approx 10 \sigma_z$, we find $\tau_d \approx 1572$ h. The depolarization arising from this source is likely to be insignificant, even with substantially more pessimistic assumptions.

### 4.2.4 Depolarisation from beam-beam forces

The fields of colliding beams can excite spin resonances associated with betatron motion. In the vertical plane these take on the form

$$\gamma_a = kp + k_z Q_z,$$  \hspace{1cm} (III.15)

where $p$, an integer, is the superperiodicity of the interaction regions. This effect is characterized by the perturbing force being periodic in revolution frequency and in the form of an abrupt localized kick.

According to Kondratenko\(^{24}\), the depolarization lifetime due to beam-beam interactions scales similarly to the normal beam-beam lifetime from non-linear resonances, since the same resonance structure is involved in both cases. Thus it is claimed that operation comfortably below the beam-beam limit will ensure an adequately long depolarization time from this effect.

In CHEEP the electron linear tune shift from the proton beam is less than 0.02. An approximate evaluation of Kondratenko's criteria indicates a depolarization time in the range of hundreds of hours. This results largely from the small vertical beam size expected in CHEEP, and therefore all efforts should be made to keep it small.

### 4.2.5 Vertical dispersion

Although the depolarizing action of vertical dispersion can be fully taken into account in the general formalism of Derbenev and Kondratenko\(^{25}\), we have as yet no explicit approximate formula for this effect. However, from their derivation of the formula for closed-orbit perturbations [our Eq. (III.7)], it is possible to deduce the form of the contribution to $\tau_d^{-1}$ arising from the variation of vertical closed orbit with momentum.

On the assumption that the Fourier spectrum of vertical dispersion is similar to that of vertical closed-orbit distortions, we conclude that the contributions to $\tau_d^{-1}$ from the two effects are approximately equal if

$$\tilde{\delta}_z \approx 2(\gamma_a)^2,$$

where $\tilde{\delta}_z$ is the peak vertical dispersion and $2$ the peak c.o. amplitude. Thus in CHEEP, $\tilde{\delta}_z$ should be kept below about 0.2 m to avoid undue depolarization from this source.

### 4.2.6 Reduction of polarization caused by special magnets

The magnet configuration around the interaction region is determined by collision-geometry requirements and by the need to rotate the polarization vector parallel to the electron beam. As a result, the polarization vector here does not, in general, lie parallel to the magnetic field as in the normal lattice magnets. Thus the radiative polarization effect of quantum emission in these magnets acts in a different effective direction from that of the normal lattice and tends to reduce the asymptotic value of polarization in the required direction.
It should be noted that this effect is quite distinct from the depolarizing effects discussed above, since it is not associated with spin resonances or with the damping of trajectory oscillations.

Following Derbenev, Kondratenko and Skrinskii\textsuperscript{23)} the asymptotic degree of polarization can be written as

\[
p_\infty = \frac{8}{5\sqrt{3}} \frac{\langle \hat{n} \cdot \vec{B} \rangle}{|\rho|^3} \left( 1 - \frac{2}{3}(\hat{n} \cdot \vec{v})^2 \right) \leq \frac{8}{5\sqrt{3}} \tag{III.16}
\]

where \(\rho\) is the radius of curvature and \(\hat{n}, \vec{B},\) and \(\vec{v}\) are unit vectors in the directions of the principal polarization eigenvector, the magnetic field, and the velocity, respectively. The averages are taken around one revolution of the machine. Evaluating Eq. (III.16) for the current CHEEP parameters shows that, with a single interaction region, the effective polarization is reduced below the ideal \(8/5\sqrt{3}\) by about 12.5\%, which is uncomfortably large. With two similar interaction regions of this kind, the loss of polarization would hardly be acceptable.

In fact, about 60\% of this lost polarization results from the vertical bends required to restore the electron orbit to the correct position in the SPS tunnel. A reduction in the vertical separation between the electron ring and the SPS would most likely permit weaker vertical bending in these regions and a useful improvement in the maximum polarization. Furthermore, there are in principle other magnet configurations which achieve the required spin rotation with suitable geometry, and which may be more favourable in this context.

4.3 Special measures

Having taken all precautions in the general design of CHEEP to obtain a high level of intrinsic transverse polarization, special measures are required to correct for errors, measure the degree of polarization, and maximize the utility for e-p physics. Some of these topics are briefly discussed below.

4.3.1 Longitudinal polarization, both helicities

The present arrangement of the CHEEP interaction region\textsuperscript{15)} uses a variable geometry of vertical and horizontal bends to obtain a longitudinal polarization at the interaction point over a wide energy range. With this layout the natural helicity is negative for both electrons and positrons; the addition of the opposite helicities would be a valuable asset for the physics programme.

We know of four possible ways of reversing the direction of the polarization vector at the interaction point, namely:

a) spin reversal by crossing a resonance;

b) variable geometry near the interaction region\textsuperscript{15)};

c) "reverse-kink" wigglers\textsuperscript{22)};

d) two electron channels near the interaction region.
Of these, (a) is potentially the most simple and attractive, (b) can be achieved over a limited energy range near 25 GeV with the present layout, (c) is not favoured because of synchrotron-radiation problems, and (d) would be very clumsy and more costly.

A preliminary study of the spin-reversal method (a) has been made by Montague. The results suggest that a fairly high degree of spin reversal might be obtained by a carefully controlled resonance crossing; however, the criteria for adiabaticity and spin correlation against quantum fluctuations are not sufficiently strongly satisfied to permit the efficiency of the process to be estimated at present, and more detailed studies will be necessary.

The spin-reversal technique has the advantage that the change in helicity can be made without disturbance to the interaction region geometry and the detector calibration. It has the slight drawback that the reverse polarization decays with the time constant $\tau_p$ and has to be "refreshed" at intervals with another resonance crossing and activation of the "kink" magnets.

### 4.3.2 Vertical closed-orbit control

It is evident from Section 4.2.2 that CHEEP requires a degree of vertical closed-orbit control going somewhat beyond the conventional needs in terms of aperture. However, the conditions are likely to be only slightly more stringent than those currently achieved in the SPS, where peak c.o. errors of $\sim 2.5$ mm have been measured (L. Burnod, private communication). The Fourier analysis of the orbit yields harmonic coefficients in quite good agreement with the Lorentz spectrum of Eq. (11.12) up to the limit of system resolution ($k = 5^4$).

In CHEEP we propose to install beam pick-ups at every quadrupole, 216 in all, which should provide enough resolution for harmonic numbers up to at least $k \sim 70$. This covers the range of $(\gamma a)$ up to 30 GeV and leaves a reserve for the nominal energy of 25 GeV where $(\gamma a) \sim 57$.

It would of course be rash to assume that the amplitudes of such high harmonics could be extracted from a Fourier analysis with the necessary accuracy and confidence, even with the reserve of measuring points foreseen. The purpose of such a system is to ensure, firstly, that a measurable degree of polarization can be achieved before having to measure it and, secondly, that changes in the relevant region of the harmonic spectrum can be checked during orbit-correction procedures.

Once conditions have been established to permit a measurable degree of polarization to develop, systematic corrections to two or three harmonics near $(\gamma a)$ can be made, using the polarization level as a measure of success.

### 4.3.3 Determination of the transverse polarization

To permit setting up of polarized beams in CHEEP, a fast polarimeter is essential. We propose to monitor the transverse polarization of the electrons in CHEEP by scattering circularly polarized photons on the circulating electron beam. The cross-section for this process is given by

$$\frac{d^2\sigma}{d\Omega d\phi} = \frac{d^2\sigma_2}{d\Omega d\phi} \cdot P \cdot \frac{d^2\sigma_2}{d\Omega d\phi} \cos \phi.$$
The various quantities are defined as follows:

\[ \phi = \text{angle between the scattering plane and the plane defined by the polarization vector and the momentum of the electron;} \]

\[ \rho = \frac{k}{k_{\text{max}}}, \text{ } k \text{ is the energy of the back-scattered photon and } k_{\text{max}} \text{ its largest value;} \]

\[ k_{\text{max}} = 4a \gamma^2 k_i; \]

\[ k_i = \text{energy of the incident photon;} \]

\[ \gamma = \text{electron Lorentz factor;} \]

\[ a = \frac{1}{1 + 4\gamma k_i/m_e c^2}; \]

\[ P_e \text{ and } P_\gamma \text{ are the magnitudes of the transverse polarization of the electron and the circular polarization of the photon, respectively.} \]

The first term in the cross-section has no \( \phi \)-dependence and is given by

\[ \frac{d^2\sigma_1}{d\phi d\phi} = r_8^2 a \left[ \frac{\rho^2(1 - a)^2}{1 - \rho(1 - a)} + 1 + \left[ \frac{1 - \rho(1 + a)}{1 - \rho(1 - a)} \right]^2 \right] \]

with \( r_8 = 2.81 \times 10^{-13} \text{ cm.} \)

The second differential cross-section is given by

\[ \frac{d^2\sigma_2}{d\phi d\phi} = r_8^2 a \left[ \frac{\rho(1 - a) \sqrt{\lambda a \rho(1 - \rho)}}{1 - \rho(1 - a)} \right]. \]

\( P_e \) can then in principle be determined from a measurement of the cross-section at \( \phi = 0^\circ \) and \( 180^\circ \) assuming \( P_\gamma \) to be known. The cross-section \( d^2\sigma_2/d\phi d\phi \) and the asymmetry ratio

\[ A = \frac{d^2\sigma_2/d\phi d\phi}{d^2\sigma_1/d\phi d\phi} \]

are plotted in Fig. III.23 for \( k_i = 3.5 \text{ eV} \) and \( E_B = 25 \text{ GeV} \) (\( a = 0.42728 \)) as a function of \( \rho \). Note that the asymmetries are rather large, reaching values around 0.3 at \( \rho = 0.7 \). The scattered photons are in a cone with an opening angle \( \theta \) with respect to the electron direction. There is a one-to-one correspondence between the scattering angle and the photon energy, i.e.

\[ \theta = \frac{1}{\gamma} \left[ \frac{1 - \rho}{\sqrt{\lambda a \rho}} \right]^{\frac{1}{2}}. \]

The values for \( \theta \) and \( k \) are listed below in Table III.4 for \( k_i = 3.5 \text{ eV} \) and \( E_B = 25 \text{ GeV.} \).
Table III.4
Parameters of back-scattered photons

<table>
<thead>
<tr>
<th>ρ</th>
<th>$\theta$ (10^{-5} rad)</th>
<th>k (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.1</td>
<td>9.39</td>
<td>1.44</td>
</tr>
<tr>
<td>0.2</td>
<td>6.26</td>
<td>2.87</td>
</tr>
<tr>
<td>0.3</td>
<td>4.78</td>
<td>4.31</td>
</tr>
<tr>
<td>0.4</td>
<td>3.83</td>
<td>5.74</td>
</tr>
<tr>
<td>0.5</td>
<td>3.13</td>
<td>7.18</td>
</tr>
<tr>
<td>0.6</td>
<td>2.56</td>
<td>8.62</td>
</tr>
<tr>
<td>0.7</td>
<td>2.05</td>
<td>10.05</td>
</tr>
<tr>
<td>0.8</td>
<td>1.57</td>
<td>11.49</td>
</tr>
<tr>
<td>0.9</td>
<td>1.04</td>
<td>12.92</td>
</tr>
<tr>
<td>1.0</td>
<td>0</td>
<td>14.36</td>
</tr>
</tbody>
</table>

The scattering angles are rather small, and a special high-beta section is needed to obtain a nearly parallel beam with divergence angles smaller than the scattering angle. The electrons and the back-scattered photons are separated by bending magnets, and the photon detectors -- in order to define $\varphi$ -- must be mounted of the order of 100 m downstream of the interaction point.

The measurement can in principle be made as follows: linearly polarized light from a laser is passed through a $\lambda/4$ polarizer and turned into right (left) circularly polarized light. The relative light output of the laser is monitored by a diode. The laser is mounted above the beam plane and the light, by means of mirrors, is made to collide nearly head-on with the circulating electron beam. The back-scattered photons are separated from the electrons by passing the beam through bending magnets, and measured using NaI crystals. In principle the azimuthal angles can be defined by a set of slits; however, the small scattering angles and the finite beam divergences make this very difficult. We therefore propose to use the following scheme: a coarse collimator is used to define a $\varphi$ range between 20° and 160° or between 200° and 340°, respectively, for $0.8 \geq \rho \geq 0.2$. The exact scattering angle is then defined by a measurement of the photon energy. The asymmetry is determined by observing the different rates for the two $\varphi$ ranges. Note that the asymmetry can also be measured for a fixed $\Delta \varphi$ range by flipping the helicity of the incoming photon. Next we estimate the rates. The luminosity for $\gamma$-e collisions is given by

$$L = 2 \frac{\mathcal{N}\mathcal{E}}{\mathcal{S}} \sqrt{\frac{L}{\gamma \mathcal{C} f_0}}.$$

Here $\mathcal{N}\mathcal{E} = 1.5 \times 10^{13}$ is the number of circulating electrons; $\mathcal{N}\gamma$ the number of photons per second; $\mathcal{S}$ and $L$ are the cross-section and the length of the interaction region, respectively; $f_0 = 45.38$ KHz is the revolution frequency; and $c = 3 \times 10^{10}$ cm/sec the velocity of light. We assume a laser with an instantaneous power of 1000 W and an average power of 0.1 W. This yields a peak flux of $1.79 \times 10^{21}$ photons/sec or an average flux of $1.79 \times 10^{17}$ photons/sec.
To evaluate $S$ we assume $\beta = 500$ m in both planes. From the normalized emittance $\sigma_x^2/\beta = 7.7 \times 10^{-8}$ rad m and $\sigma_z^2/\beta = 0.35 \times 10^{-8}$ rad m we find $\sigma_x = 0.62$ cm, $\gamma_x = 1.2 \times 10^{-5}$ rad, $\sigma_z = 0.13$ cm, and $\gamma_z = 2.6 \times 10^{-6}$ rad. This yields $S = 4\pi \sigma_x \sigma_z = 1.05$ cm$^2$. Also note that the vertical divergence angles $\gamma_z$ are small compared to the scattering angles. We further assume that the photon and electron beams collide over a distance of 10 m. The $y\gamma$-cross-section integrated over the azimuthal acceptance for $\rho < 0.8$ yields $4.6 \times 10^{-26}$ cm$^2$, leading to a counting rate of about 300 events/sec. The background caused by electron bremsstrahlung on the residual gas was evaluated assuming a pressure of $10^{-9}$ Torr (mainly $H_2$). This background was found to amount to only a few percent of the signal. Another source of background is the synchrotron radiation caused by the dipole magnets and the quadrupoles upstream of the interaction region. This background is presumably not serious; however, it cannot be evaluated before the high-beta section has been designed.

5. **Electron Ring**

5.1 **Lattice**

The electron lattice is derived from the SPS lattice with one quadrupole and four bending magnets per half cell, corresponding elements being superimposed. The electron ring is supported 1.1 m above the proton ring (centre-to-centre), which allows easy access and does not disturb the survey system. The phase advance per cell of $\pi/2$ is maintained and the quadrupole strengths near the long straight sections are adjusted to obtain regular betatron and dispersion functions in the arcs (unlike the proton lattice where the dispersion is allowed to vary considerably around the circumference). The phase advance is higher than most other electron storage rings but it has been shown that chromatic defects are correctable in this type of lattice\(^{31}\) and, in fact, this structure has previously been proposed for an electron ring\(^{32}\). Imposing regular dispersion in the bending arcs results in the smallest beam size for a machine with this lattice and therefore improves the luminosity. The aperture in the ring has been calculated for $\pm 10\sigma$ and adding $\pm 2.5$ mm for closed orbit (as measured in the SPS). As is normal in electron rings, the uncoupled horizontal emittance and the fully coupled vertical emittance are used. The requirement at 30 GeV is remarkably similar to that in PETRA and it was decided to adopt for CHEEP exactly the same vacuum chamber as for PETRA. Many of the engineering techniques already developed at DESY are therefore directly applicable and the design of other components is greatly facilitated.

The "standard" straight sections of the SPS tunnel (i.e. those not earmarked for future use as e-p experimental areas) contain the proton injection (LSS 1), ejection (LSS 2, LSS 6) and RF (LSS 3), and the electron ring should not impede access to the SPS equipment. Considerable improvement is obtained if the electron ring is displaced outwards in these regions by altering the position of the missing bends situated either side of all straight sections. A convenient choice is to leave out the two bending magnets exactly one half cell away, giving a displacement of 0.56 m. This provides the maximum displacement in the straight section consistent with the limited space available in the arcs of the tunnel.

The characteristics of the electron beam are solely determined by the guide field containing it\(^{33}\), and the design of the lattice must provide the required characteristics over a range of energies. In a ring with fixed geometry this is not possible and the designer must resort to some sort of variable geometry in which the average bending radius is altered.
to try to counteract the variation of the beam characteristics with energy. In CHEEP the polarization of the electron beam is considered to be of prime importance and this precludes the use of reverse bends in any such scheme. A new method is proposed for CHEEP\textsuperscript{22} in which a short high-field magnet is placed between the standard length bending magnets in an otherwise normal cell. Then the required bending per cell may be obtained by a mixture of the short strong bend ("kink") and the weakened standard bends. The limit is then given when the weakened lattice bends are switched off completely with the kink magnet at its maximum field. The most suitable position for a kink magnet in the CHEEP lattice is between the two bending magnets nearest to the horizontally defocusing (D) quadrupole. These two magnets and the kink magnet then make up the variable geometry unit. This arrangement has two advantages. Firstly, the vertical beam size is large at the kink magnet, which reduces the synchrotron radiation intensity on the vacuum chamber walls. Secondly, the horizontal beam size is small, which offsets the effect of sagitta variations and enables the standard vacuum chamber to be used (Fig. III.24). The range of energies over which optimization is required is 15-25 GeV. A value of 0.4 m for the length of the kink magnet has been adopted and implies a maximum field of 2.1 tesla.

Approximate expressions for the effect of the kinks on the electron beam characteristics have been derived\textsuperscript{22}. The following symbols are used:

- \( E \) - electron energy in GeV
- \( N_k \) - number of kinks in the ring
- \( L_k \) - length of each kink magnet: 0.4 m
- \( N_m \) - total number of standard bending magnets: 744
- \( L_m \) - length of each standard bending magnet: 6.26 m
- \( f_0 \) - revolution frequency: 43.4 kHz
- \( L \) - circumference of electron ring: 6911.5 m
- \( \rho_0 \) - radius of curvature in standard bending magnets: 741.25 m
- \( \rho_k \) - radius of curvature in kink magnets
- \( J_x \) - damping partition coefficient: \( J_x \approx 1, J_e \approx 2 \).

Then

\[
\text{Polarization time } \tau_{pol} \approx \frac{98L}{E^5} \frac{\rho_0^3}{N_m L_m} \left[ 1 + \frac{N_k L_k}{N_m L_m} \left( \frac{\rho_k}{\rho_0} \right)^3 \right]^{-1} \text{ sec},
\]

\[
\text{Synchrotron energy loss per turn } U_{tot} \approx \frac{8.85 \times 10^{-2} E^4}{\rho_0} \left[ 1 + \frac{N_k L_k}{N_m L_m} \left( \frac{\rho_k}{\rho_0} \right)^2 \right] \text{ MeV},
\]

\[
\text{Damping rate } \alpha_{\text{damp}} \approx 4.425 \times 10^{-5} \frac{J f_0 E^3}{\rho_0} \left[ 1 + \frac{N_k L_k}{N_m L_m} \left( \frac{\rho_k}{\rho_0} \right)^2 \right] \text{ sec}^{-1},
\]

\[
\text{Energy spread } \sigma_E/E \approx 1.212 \times 10^{-3} \sqrt{\frac{1}{J_e \rho_0} \left[ 1 + \left( \frac{N_k L_k}{N_m L_m} \left( \frac{\rho_0}{\rho_k} \right)^3 \right) \right] \left[ 1 + \left( \frac{N_k L_k}{N_m L_m} \left( \frac{\rho_0}{\rho_k} \right)^2 \right) \right]^{1/2}},
\]
Horizontal emittance \( \varepsilon_X = \frac{1.47 \times 10^{-6} E^2}{J X \rho_0} \left[ \frac{1 + \left( N_k \beta_k / N_m \beta_m \right) (\rho_0/\rho_k)^3 (W_k/W_0)}{1 + \left( N_k \beta_k / N_m \beta_m \right) (\rho_0/\rho_k)^2} \right] \),

where the Courant and Snyder invariant \( W = \frac{D_x^2 + (D_x / X)^2}{\beta_x X} \).

It is now necessary to choose the value of \( N_k \) and the mode of operation. At top energy the natural polarization time is 100 minutes, and this should be reduced. This can be achieved with the minimum extra RF by using a few kinks at maximum field rather than all the kinks at a lower field. At slightly lower energies the available RF will be sufficient to run all of the kinks installed even at the maximum field if so required. The effect of the kinks is to reduce the polarization time (good), reduce the damping times (good), increase the energy spread (bad), and modify the emittance (depending on the value of \( W_k/W_0 \)). The preferred mode of operation is to use, at any energy, the maximum number of kinks, and to vary the field in all of them to obtain the desired polarization time. This minimizes the increase in energy spread. The total number of kinks in the ring is dictated by the polarization time at the lowest energy, and 16 kinks would be needed to obtain about 30 minutes using the maximum field. At top energy only 4 kinks will be used (implying two sets of kinks independently powered). The \( W_k/W_0 \) ratio must be less than 0.2 to avoid a large increase in beam size at low energy. This means a region of low dispersion, which can be obtained by incorporating the kinks into the matching regions either side of the standard straight sections. The kinks are placed in pairs either side of the D-quadrupole with the lowest dispersion in each of the eight matching sections. This now freezes the geometry of these regions so betatron matching can be attempted, and one solution is shown in Fig. III.25. This structure was then used for computer calculations of the beam characteristics including the effect of the kinks, and good agreement with the analytic formulae was found. However, an unexpected result was that at low energies the extra aperture required to accommodate the increased energy spread outweighed any reduction due to emittance shrinkage. This imposes an additional limitation on the maximum field that can be used in the kink magnets and hence limits the improvement in polarization time. Figure III.26 shows the variation of polarization time actually attainable with the system. The maximum time of 45 minutes still seems acceptable. The synchrotron radiation loss per turn is shown in Fig. III.27. The increase is only 11% at 25 GeV (this implies that the maximum energy without kinks is 25.66 GeV). The drop at 24.6 GeV due to changeover from 4 to 16 kinks is 5 MeV (6%). The damping times remain relatively constant in the range 15-25 GeV but increase rapidly at lower energies, reaching 1.34 seconds at 5 GeV. It therefore appears that wigglers will be required at injection\(^{3\text{b}}\) and these would be placed in the missing magnet regions (Fig. III.25). The relative energy spread is shown in Fig. III.28. This is rather large and implies that chromaticity correction is required over a large momentum bite. This has not yet been attempted.

5.2 Main magnet system

5.2.1 Introduction

The provisional design of the main magnet system has been greatly influenced by the particular requirements of the installation in the SPS ring, i.e.
radial and vertical dimensions as small as possible,
- weight as low as possible,
- limited power consumption in order to take advantage of the SPS cooling plant,
- limited space for magnet supports.

To take advantage of the existing distribution of cooling water in the SPS, the coils of the CHEEP magnets will be cooled in parallel with the copper coils of the SPS magnets. Thus we had to choose copper for the coils of CHEEP to avoid corrosion problems. This choice is also favourable for the quadrupole design, because of the constraints on the overall magnet dimensions.

In the following a description of the dipole and quadrupole magnets will be given. Their main parameters are summarized in the general parameter list.

### 6.2.2 Dipoles

#### Cores

Coil designs with bent ends were abandoned as they are not necessary for the required field precision and would increase the total magnet cost by 30%. The choice is restricted to H-magnets (Fig. III.29) with good mechanical stability and field homogeneity, or C-magnets (Fig. III.30) which allow an easy access to the vacuum chamber and the installation of vacuum pumps with a lateral connection to the vacuum pipe.

Detailed computation\(^{15,16}\) of the magnetic field shows that shimming of the pole edges allows a decrease in the pole width to 180 mm. The field homogeneity for this pole configuration is found to be:

\[
\Delta B/B = 3.5 \times 10^{-4} \text{ for the H-magnet (gap height = 72 mm, Fig. III.31) ,}
\]
\[
\Delta B/B = 7.6 \times 10^{-5} \text{ for the C-magnet (gap height = 77 mm, Fig. III.32) .}
\]

The saturation of the steel can be expressed by

\[
\Delta B/B = \left[ B(E) - B'(E) \right] / B'(E) ,
\]

where \( B(E) \) is the central field at energy \( E \), and \( B'(E) \) is the field calculated at the same current for an infinite permeability; \( \Delta B/B \) is equal to \( -3.8 \times 10^{-3} \) at 50 GeV/c for the H-magnet (for the C-magnet it is \( \Delta B/B = -5.8 \times 10^{-3} \)).

The cores are assembled from steel laminations held together by longitudinal flat straps which are welded on the laminations and to the end plates. Mixing laminations according to their measured magnetic properties will reduce the difference of the bending power between dipoles. Moreover, the spread in bending power of the dipoles will be reduced to about \( 10^{-4} \) by addition or subtraction of shims to the core ends.

The magnetic forces acting on the poles may be neglected; the longitudinal mechanical stability is, however, more critical owing to the core length of 6 m and the small cross-section. A reinforcing structure is used to increase the moment of inertia of the assembly. For the H-magnet, a box-shaped beam containing the magnet has been chosen because of its favourable ratio of moment of inertia to over-all weight; the magnet is supported inside the beam by means of screws, which allow correction of a possible twist of the magnet core. For the C-magnet, this box would be placed above the core. The beam weight may be reduced
by adopting a variable cross-section or by supporting it at the two points giving the minimum deflection. As an example, a 260 kg beam loaded with the 3 tons of the 6 m long H-magnet will have a maximum deflection of 0.3 mm with optimal two-point support.

**Coils**

The power dissipation in the dipole coils has been limited to 3.3 MW, a total of about 5 MW being foreseen for the main magnet system (dipoles, quadrupoles, busbars).

The 6.095 m core length corresponds to a magnetic length of 6.202 m (H) or 6.214 m (C) (taking into account the average thickness of the end shims) for which a central field of 0.1363 T (H) or 0.1360 T (C) is required at 30 GeV/c.

This field may be obtained by a coil having three turns per pole and a current of 1306 A (H) or 1397 A (C); the conductor has a cross-section of 55 mm x 18 mm. The voltage drop across the dipoles will be 2.52 kV. A 4 kV power supply, centre-tapped for earth, will be required for exciting the dipoles and quadrupoles connected in series by appropriate busbars. The higher current in the C-magnet is due to the gap height of 77 mm. For a current of 1306 A, as for the H-magnet, the C-core length should be 6.261 mm.

One water-cooling circuit per magnet is foreseen.

**5.2.3 Quadrupoles**

The quadrupole design is determined by the following requirements:
- Use of the PETRA vacuum chamber with its distributed vacuum pump.
- Possibility of electrical series connection with the dipoles, to save on busbar cost. This limits the number of ampere-turns for the quadrupole to an integer times the dipole current.
- Quadrupole length of about 3 m as in the SPS; a reduction of this figure would reduce the beam-damping aperture.
- Quadrupole maximum half height of 250 mm, because of the separation between the electron and proton beams (1100 mm) and the free space required for the alignment equipment for both rings.

According to beam optics calculations, the integrated gradient at 30 GeV/c must be 4.5675 T, at the nominal machine working point $Q = 27.5$. Differences in excitation between the focusing ($Q_\uparrow$) and defocusing ($Q_\downarrow$) quadrupoles will be obtained by means of auxiliary windings around the poles. These windings will have a capability of ±10% of the maximum gradient.

Preliminary AGS computations, using the SPS lattice, have shown that a tuning range of ±2.5 in $Q$ may be achieved by these auxiliary windings. Low-current busbars (about 200 A), easy to install in the SPS tunnel, will be used for their electrical connections.

**Quadrupole layout**

A fully symmetric quadrupole was chosen in order to achieve a high purity of the field. The relationship between integrated gradient ($G_\perp$), ampère-turns (NI), magnetic length ($\ell_m$), and half aperture (r) is given by:

$$r^2 \approx 2 \times NI \times \mu_0 \times \ell_m / G_\perp,$$

(III.17)

where $G_\perp$ is fixed, $\ell_m$ may vary between 2.8 and 3 m, NI = nLI dip, with n = integer.
Once \( r \) has been chosen, the steel length will be adjusted to obtain the required magnetic length; this adjustment will be of the order of a few percent.

Assuming \( I_{\text{dip}} = 1306 \text{ A} \) (which is the nominal current for the \( \pi \)-type dipoles at 30 GeV/c), we see from formula (III.17) that only \( n = 3 \) and \( n = 4 \) can be accepted.

For \( n = 3 \), an aperture of 160 mm at \( x_m = 3 \text{ m} \) is needed to obtain the required strength. A magnet layout of this type is shown in Fig. III.33; its parameters are those shown in the main parameter list.

The field distribution in the pump channel is shown in Fig. III.34. Because of the space taken by the distributed vacuum pump, a pole width of 80 mm is about the maximum acceptable in view of the tolerances on the vacuum chamber dimensions and of the thickness of lead shielding against synchrotron radiation.

Several MAREA computations\(^{16} \) were run to optimize the pole profile; the results are shown in Fig. III.35. One can see that a gradient homogeneity \( \Delta G/G \) of about \( \pm 0.6\% \) may be achieved in the aperture necessary for the beam. \( \Delta G/G = \pm 0.15\% \) could be obtained by increasing the pole width to 96 mm at the expense of the lead shield and of a modification of the pump channel.

For \( n = 4 \), an aperture of 180 mm at \( x_m = 2.8 \text{ m} \) is needed. The larger aperture permits us to choose a wider pole (e.g. 100 mm), which in turn could make it possible to achieve, according to MAREA calculations, \( \Delta G/G < 0.1\% \) across the whole aperture needed at 30 GeV/c. However, this would be at the price of an increase in magnet weight and power dissipation. The magnet's half height would still be below 250 mm.

Should it prove to be possible to omit the slot for the distributed vacuum pump inside the quadrupole, a magnet design with a bore of about 100 mm, such as that chosen for PETRA and PEP, will be studied. Such a design offers a saving in power dissipation of about 500 kW with respect to the one shown in Fig. III.35. However, it would require a construction in which the core is split into four quadrants, making it more difficult and expensive to assemble.

Core and coil design

Since we have a low-field magnet, a pole with parallel sides is used; this leads to very simple race-track coils and is also favourable from the weight point of view. At the expense of a small increase in power dissipation, the coil is designed (see Fig. III.33) to be slid on to the poles without splitting the magnetic circuit into four quadrants. The magnet laminations may be punched in one piece, thus making the core assembly easier than in the case of core quadrants.

Reinforcing steel straps are welded to the sides of the quadrupole, in order to keep its deflection below 0.1 mm. Their shape and thickness will depend on the choice of the magnet supports.

5.3 The electron RF system

5.3.1 Introduction

The SPS RF system for protons is not equally suitable for acceleration of the electron beam in CHEEP. The proton RF system is designed for uniform beam loading, a relatively large bandwidth, and a modest voltage gradient, resulting in a cavity structure of rather low impedance.
In contrast, the electron bunches in CHEEP fill only a small number of the available RF buckets, resulting in a heavy transient beam loading. Furthermore, much higher voltages are required to make up the energy losses due to synchrotron radiation and beam-cavity interactions, and to provide the required overvoltage factor for quantum lifetime. Such high voltages in turn require cavities of high shunt impedance to minimize the RF power loss in the cavities themselves.

Because of the high RF power required, the use of high-efficiency klystrons is preferred.

Thus there is not very much 200 MHz equipment that could be copied from existing SPS equipment. Furthermore, all the low-level RF is at an intermediate frequency of 10.7 MHz; for a frequency different from 200 MHz, only preselection and oscillator frequency need to be readjusted, provided one uses the same low-level RF system in the e-ring.

It seems that no special argument exists in favour of adopting the same frequency as in the proton ring. The harmonic numbers of the electron and the proton ring do not need to be equal; the frequencies must simply be multiples of the bunch number.

6.3.2 Choice of frequency

If we are not really confined to the SPS proton ring frequency, we have to investigate which range of frequencies could be used and whether there is an advantage in using a different one. A lower limit is immediately apparent; namely, a study has shown that RF cavities, larger in diameter than the existing 200 MHz cavities for the protons, cannot be installed because of lack of space. This means essentially that the lowest possible frequency tends to be even higher than 200 MHz because of the more voluminous high-Q cavities. The upper frequency limit is given by the maximum tolerable $Q_S$. If we wish to stay below a $Q_S = 0.1$, the maximum possible frequency for the final energy level of 30 GeV will be around 400 MHz.

In Table III.5 the most important data for this frequency range are listed both for the design energies of 25 GeV and 30 GeV.

This table is based on the following parameters:

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
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</tr>
<tr>
<td>Bending radius</td>
<td>741 m</td>
</tr>
<tr>
<td>Momentum compaction factor</td>
<td>$1.66 \times 10^{-3}$</td>
</tr>
<tr>
<td>Number of electrons</td>
<td>$1.5 \times 10^{13}$</td>
</tr>
<tr>
<td>Quantum lifetime</td>
<td>100 h</td>
</tr>
<tr>
<td>Shunt impedance of RF cavities at 200 MHz</td>
<td>20 MΩ/m</td>
</tr>
<tr>
<td>&quot; &quot; &quot; &quot; 400 MHz</td>
<td>28.3 MΩ/m</td>
</tr>
<tr>
<td>Total RF cavity length at 25 GeV</td>
<td>96 m</td>
</tr>
<tr>
<td>&quot; &quot; &quot; &quot; 30 GeV</td>
<td>192 m</td>
</tr>
</tbody>
</table>

There are no models of a possible e-ring RF cavity available yet for measuring the precise shunt impedance. Instead, we have assumed values which might be expected from scaling of measurements on side-coupled or similar cavities (Los Alamos, Chalk River).

The cavity length is a result of layout studies for the SPS tunnel and, according to a certain cost optimization, involves the minimum of capital costs.
The somewhat surprising result in comparing the data of Table III.5 is that nothing really seems to favour the higher frequency because even the slightly lower RF losses will probably be counterbalanced by more severe losses from higher-order modes because of the shorter bunches. In practice, it may be that equipment for lower frequencies is more expensive. However, preliminary investigations with industry have shown that the cost variation is rather small within our frequency range. On the other hand, the scaling of the cavity impedance should be more favourable for the lower frequency because in reality the hole for the beam will be relatively smaller. As the advantage of a longer bunch and a lower $Q_s$ is with the lower frequency, we choose here the lowest possible frequency for which we can fit a cavity of the Chalk River type into the tunnel. A 3 GHz model scaled to an inside diameter of ≤ 75 cm and to a beam hole of = 10 cm diameter yields a frequency of = 500 MHz.

The detailed choice of the harmonic number was made after examination of possible options to reduce the positron filling time, as will be explained in Section 8. One of these options is to put the required number of particles into more than 60 buckets at injection, which takes less time than filling the particles directly into 60 buckets. After filling, the bunches are combined via a by-pass such that only the 60 master buckets are filled in the end. In order to have maximum latitude for such a bunch accretion scheme, a harmonic number of 7200 is chosen which, in principle, permits initial distribution of the particles over 60$n$ (n = 2, 3, 4, 5, 6) buckets, where $n$ is also the reduction factor of the positron filling time. Thus the frequency becomes 512.3 MHz, which appears to be a rather safe compromise. The relevant data are given in Table III.5.

### Table III.5

<table>
<thead>
<tr>
<th></th>
<th>1</th>
<th>2</th>
<th>3</th>
<th>4</th>
<th>5</th>
<th>6</th>
</tr>
</thead>
<tbody>
<tr>
<td>Kinetic energy</td>
<td>25.00</td>
<td>25.00</td>
<td>25.00</td>
<td>30.00</td>
<td>30.00</td>
<td>30.00</td>
</tr>
<tr>
<td>Frequency</td>
<td>200.40</td>
<td>312.31</td>
<td>400.79</td>
<td>200.40</td>
<td>312.31</td>
<td>400.79</td>
</tr>
<tr>
<td>Overvoltage factor</td>
<td>1.35</td>
<td>1.48</td>
<td>1.58</td>
<td>1.31</td>
<td>1.43</td>
<td>1.51</td>
</tr>
<tr>
<td>RF voltage</td>
<td>62.82</td>
<td>69.0</td>
<td>73.57</td>
<td>126.72</td>
<td>138.1</td>
<td>146.44</td>
</tr>
<tr>
<td>RF voltage gradient</td>
<td>0.65</td>
<td>0.72</td>
<td>0.77</td>
<td>0.66</td>
<td>0.72</td>
<td>0.76</td>
</tr>
<tr>
<td>Losses in lattice</td>
<td>4.86</td>
<td>4.86</td>
<td>4.86</td>
<td>10.08</td>
<td>10.08</td>
<td>10.08</td>
</tr>
<tr>
<td>Losses in cavities</td>
<td>2.05</td>
<td>1.98</td>
<td>2.00</td>
<td>4.18</td>
<td>3.97</td>
<td>3.95</td>
</tr>
<tr>
<td>Stable phase angle</td>
<td>48.02</td>
<td>42.6</td>
<td>39.40</td>
<td>49.83</td>
<td>44.53</td>
<td>41.39</td>
</tr>
<tr>
<td>Synchrotron frequency</td>
<td>1.96</td>
<td>2.7</td>
<td>3.23</td>
<td>2.50</td>
<td>3.43</td>
<td>4.10</td>
</tr>
<tr>
<td>$Q_s$</td>
<td>0.045</td>
<td>0.062</td>
<td>0.075</td>
<td>0.058</td>
<td>0.079</td>
<td>0.095</td>
</tr>
<tr>
<td>Minimum bunch length</td>
<td>0.106</td>
<td>0.077</td>
<td>0.064</td>
<td>0.100</td>
<td>0.073</td>
<td>0.061</td>
</tr>
</tbody>
</table>

In the final layout of the RF power plant we have to make allowance for higher-order mode losses, which will need careful design of the beam chamber and any other equipment in interaction with the beam. For the energy of 25 GeV the estimates are ~ 200 kW for higher modes and ~ 900 kW for additional radiation by the kink magnets. As a result, somewhat more than 8 MW of RF power will be required for 25 GeV and this will have to be doubled for 30 GeV.
5.3.3 Layout of the RF system

The lattice of the electron ring will be almost identical to that of the proton ring, with long straight sections in the six straight parts of the SPS tunnel (cf. Fig. III.1). The RF system and injection equipment of the proton synchrotron fill two of these straight parts; two are foreseen as interaction regions. This leaves the two straight parts where the proton ejection is installed. Since the tunnel is enlarged in these places, space is available for the RF system of the e-ring.

A non-negligible problem in this context is the fact that not only does appropriate accommodation have to be found for the cavities, but also for the high-power feeder lines, especially in the access pits. This requires a careful choice of the number of cavities and their distribution over the available space.

This number should be a power of two, i.e. 4, 8, 16, or 32 per long straight section, so that the RF power can be split or recombined much more easily by making use of hybrid circuits. It so happens that each straight section contains four semi-periods of about 28 m free length and, a little further away in the tunnel in both directions, two quarter periods of about 13 m length. Following the above arguments, this means that we have either $4 \times 24$ m or $8 \times 12$ m = 96 m available per long straight section if we allow for some space in between the cavities. In this way either two quarter periods or one half period will be left free of cavities.

It is clear that a system with an extremely long total cavity length will be more economical on the power amplifier side because the shunt impedance will be high and the voltage gradient low. Hence the cost for the tanks, assuming that the space is available, will be predominant, while for a very short total cavity length, not only the capital costs but also the running costs of the power amplifiers will be important. By comparing the estimates and prices for high-power accelerating systems it was found that there is a very flat price optimum for the installed equipment, independent of the sometimes disparate cost estimates, if the cavities fill one straight section for 25 GeV and if they fill two straight sections for 30 GeV. If the power consumption is included in this consideration the optimum length is about doubled. In this event, the price for equipment will be higher by a few percent, although the total cost over ten years will be lower. As a conclusion, the following layout seems to be possible:

<table>
<thead>
<tr>
<th>Energy</th>
<th>Total length of cavities</th>
<th>Total RF power</th>
</tr>
</thead>
<tbody>
<tr>
<td>25 GeV</td>
<td>96 m</td>
<td>7.5 MW approx. opt. for installation</td>
</tr>
<tr>
<td>25 GeV</td>
<td>192 m</td>
<td>6.2 MW approx. opt. for operation</td>
</tr>
<tr>
<td>30 GeV</td>
<td>192 m</td>
<td>15.5 MW approx. opt. for installation</td>
</tr>
<tr>
<td></td>
<td>(without higher-order mode losses or enhanced radiation).</td>
<td></td>
</tr>
</tbody>
</table>

Taking into account that the infrastructure such as buildings, etc., must be doubled for the length of 192 m and that the total number of power amplifiers would not be reduced in spite of the lower output power (because of the above-mentioned power splitting), it is probably the most practical solution in the first stage to install a total cavity length of
96 m, together with a complete power amplifier plant of about 8 MW. For the second stage of 30 GeV, the same installation could be repeated in the second long straight section. Consequently, the RF power flow is limited to ≈ 8 MW per access pit, and this anyway seems to be an upper limit given the space available in the access shafts.

5.3.4 Cavities and beam loading

Up to now, we have made all power calculations by assuming a certain shunt impedance per unit length and then choosing the total length more or less in a rather free way, neglecting the fact that beam loading is not continuous. The resultant beam loading of ≈ 65% looks very impressive, but it is very doubtful whether this efficiency can be obtained in reality if no special measures are taken. The bunches will be ≈ 75 psec long; the bunch-to-bunch distance is 380 nsec and the bucket spacing is 3.5 nsec. Hence only about 120 RF cycles are available to replace the energy, removed from the cavity by the preceding bunch, before the next bunch arrives. Thus the amplifiers will practically never see a matched load. Another problem is that it is not easy to find out exactly how much of the stored energy can be extracted during one bunch passage for different cavity types. It might turn out that a higher Q-value is required than under CW conditions. This means essentially that the power figure of 8 MW is optimistic, and we should be very well aware of the fact that this power might still increase.

At the moment we are investigating a special type of cavity array which is filled up to a certain voltage gradient, in standing wave mode. This device is similar to a π/2 side-coupled cavity, the number of cells however being determined by the bunch distance. The energy flow rate in such a cavity, i.e. the group velocity, is much slower than the beam -- namely about 10% of the beam velocity. It now seems possible to superimpose a travelling wave on the standing wave such that in between two bunch passages the travelling wave will just have enough time to replace the energy consumed by one bunch. In other words, the device resembles a storage line which collects the continuous energy from the power amplifiers and passes it on to the beam in one instant, like a travelling wave structure. However, the voltage is backed up by a standing wave. This is made possible by a complete reflection at one end of the structure, whilst the other end is matched to the voltage of both added waves, and consequently the power amplifiers continuously see a matched load despite the transient nature of beam loading. It goes without saying that all this will only work for a well-defined bunch spacing.

For 60 bunches in the SPS the bunch-to-bunch distance is 115 m and a cavity, built according to this principle, will have a length of one-half of the 115 m multiplied by the ratio of group velocity over particle velocity; this corresponds to a length of approximately 6 m. This is a very comfortable value because the sum of all cavities which will fit into the available 96 m is 16, still leaving about 1 m additional space in between the cavities, and power splitting will not present any problem.

The match of the cavity input will of course depend on the beam loading, and a continuously variable coupler will be required as the beam loading will increase during acceleration. In this way it is perhaps possible to avoid the use of ferrite isolators.
6.3.5 Power amplifiers

Very little needs to be said about power amplifiers, because by now there is plenty of experience available in this field. Within the frequency range around 300 MHz the use of klystrons is standard if high power is required. Components, such as guns, collectors, or output windows, are available and the theory is well understood; thus it no longer seems difficult to produce a klystron, for any frequency within our range, capable of delivering 500 kW and more with an efficiency above 60%. The perveance is in a range which will allow operation with collector voltages lower than 60 kV so that oil insulation can be avoided.

In order to obtain 8 MW total output power, 16 klystrons -- the same number as cavities -- are required. It seems excluded to install 16 feeder lines in one access pit and, as has been mentioned above, use should be made of power combination and splitting if the klystrons are to be installed on the surface. In this way, one waveguide can probably handle the full power.

With a klystron number equal to the cavity number it is very tempting to investigate whether the klystron can be mounted directly on each cavity, perhaps in a horizontal position on top and parallel to the accelerating cavity. The length of such a klystron will in no case be more than the 6 m of the cavity.

5.4 Vacuum system of the electron ring

The proposed design of the CHEEP vacuum system can be based largely on techniques and experience gained in the construction of the vacuum systems for SPEAR, PEP, and PETRA. Since it turned out that the aperture requirements are very close to the requirements in PETRA, it was decided to use the existing PETRA chamber profile also for CHEEP. It provides enough vertical aperture $(\pm 10 \, \text{cm})$ for operation up to 50 GeV in fully coupled mode, and the horizontal aperture would suffice even at 40 GeV in the limit of vanishing coupling. Consequently, many existing design details for vacuum chamber components, such as RF transitions for bellows, aluminium to stainless steel transitions, or pump connections, are of direct use for CHEEP.

The beam lifetime $\tau$ (h) due to bremsstrahlung on the residual gas is given by

$$\tau = 2 \times 10^{-8} \frac{X_0}{(\text{M} \cdot \text{P})}$$

where $X_0$ is the radiation length (g cm$^{-2}$), M is the molecular weight, and P the pressure (Torr) in the vacuum chamber.

When the residual gas is expressed in nitrogen equivalent pressure, we find

$$\tau (\text{h}) = 2.74 \times 10^{-8} / P \, \text{(Torr)} \ .$$

A beam lifetime of 30 hours due to vacuum, as specified for CHEEP, calls for an average nitrogen equivalent pressure of $10^{-9}$ Torr in the presence of a circulating electron beam of 100 mA. The attainment of such a low average pressure requires a careful optimization of the effective pumping speed in the system, as well as special measures for reducing the thermal outgassing and the desorption from the vacuum chamber induced by synchrotron radiation.
8.4.1 Distributed and lumped pumps for CHEEP

The low magnetic bending field, about 200 G at injection and 1.1 kG at 25 GeV, represents a serious limitation to the normal distributed sputter ion pumps operated in the bending magnets.

Since an answer to the question of whether or not a distributed ion pump can be used in CHEEP is of vital importance for the study of the vacuum system and for the detailed chamber design, laboratory work on the optimization of low-field sputter ion pumps has been started. The minimum field \( B_{\text{min}} \) (G) at which the discharge in the sputter ion pump sets in is inversely proportional to the anode cell diameter \( d \) (cm) and is given empirically by\(^{38}\)

\[
B_{\text{min}} = 600/d.
\]

As a consequence of the large pump cell diameter needed for the low field operation in CHEEP (approximately 3.6 cm diameter cells), a slightly enlarged pump channel as compared to the original PETRA profile may become necessary. Preliminary results obtained with a test pump indicate that at 25 GeV an effective pumping speed of the order of 50 \( \ell /\text{sec} \cdot \text{m}^{-1} \) with a pressure of \( 10^{-9} \) Torr should be a realistic assumption for CHEEP\(^{39}\). It is assumed that this pumping speed could be achieved in the beam pipe by pumping through holes in the wall which separate it from the pump channel. Even at the injection energy of 5 GeV this pump would provide a small but nevertheless useful pumping speed, scaling approximately as \( B^2 \).

To maintain a sufficient vacuum at injection or during periods when the main magnets are completely switched off, lumped pumps with a nominal speed of about 40 \( \ell /\text{sec} \) are provided in the inter-magnet gaps at a distance of about 6.2 m. The calculated linear pumping speed\(^{40}\) averaged over one bending magnet unit and resulting from the combined action of distributed and lumped pumps is shown in Fig. III.36. Because of the low specific conductance of the vacuum chamber of 67 \( \ell /\text{sec} \cdot \text{m}^{-1} \), and because of the large pump separation of 6.2 m, increasing the lumped pumps above 40 \( \ell /\text{sec}^{-1} \) does not result in a proportional gain in the effective linear pumping speed.

To obtain the necessary effective pumping speed in the short straight sections, distributed ion pumps will be incorporated in the vacuum chambers of the quadrupole magnets.

Figure III.37 shows a possible layout of the vacuum chamber in the inter-magnet sections, which is similar to the PEP proposal. Expansion bellows, flanges, and the lumped ion pumps are contained in a section made of stainless steel and joined to the aluminium chamber by circular stainless-steel to aluminium transitions on either side. The necessary smooth continuation of the vacuum chamber section through the bellows assembly is provided by sliding RF contacts as used in PETRA.

5.4.2 Thermal outgassing and synchrotron radiation-induced gas desorption

In view of the practical limitations of the linear pumping speed which can be realized in CHEEP with lumped and with distributed ion pumps, an assessment of the expected outgassing rates is important. On the basis of the quoted values for pumping speed and the required average pressure, the linear outgassing rate should not exceed \( 1 \times 10^{-7} \) Torr \( \ell /\text{sec} \cdot \text{m}^{-1} \). This is about one order of magnitude less than in PETRA.
Within the bending magnets the predominant part of the gas load is due to desorption induced by synchrotron radiation. Thermal desorption from the aluminium chamber has to be reduced to an acceptable value of $1 \times 10^{-8}$ Torr $\ell$ sec$^{-1}$ m$^{-1}$ by baking the chamber at approximately $120-150^\circ$C [1]. Such a mild bakeout should reduce the specific outgassing rates (predominantly hydrogen) to $4 \times 10^{-12}$ Torr $\ell$ sec$^{-1}$ cm$^{-2}$ for aluminium. The principal purpose of this bakeout would be the reduction to a negligible amount of outgassing rates for heavier gases like CO, CO$_2$, H$_2$O.

The gas desorption induced by synchrotron radiation can be estimated as increasing roughly linearly with the electron energy. Scaling from PETRA design figures, the admissible desorption rate in CHEEP at 25 GeV implies a molecular desorption yield $\eta$ of not more than $1 \times 10^{-6}$ molecules per photoelectron.

Conditioning of the vacuum chamber by prolonged running-in at maximum beam current (as done at SPEAR or in DORIS) is a very inefficient process for CHEEP because it could only be done at the low injection energy. The method proposed for achieving $\eta$ values approaching $10^{-6}$ is in situ glow-discharge cleaning. This method has been tested successfully for PETRA, and desorption yields as low as $10^{-7}$ were obtained in laboratory tests [2]. However, there still remains the uncertainty that such excellent results can be obtained on the much larger scale of 6 km circumference in CHEEP. Practical experience of the long-term behaviour of a vacuum system cleaned by this method does not at present exist.

The in situ discharge cleaning can use distributed ion pumps as the anode for the glow discharge. This procedure implies a special pump design and effective protection of the insulators against metal sputtering. The distributed ion pumps in the bending magnets and in the quadrupoles will therefore serve the double purpose of pumping and of reducing the desorption yield. Without glow-discharge cleaning, the vacuum chamber might have a desorption yield of the order of $10^{-9}$ to $10^{-8}$, even after prolonged operation and conditioning at maximum beam current.

In certain sections of vacuum chamber which cannot be cleaned by an in situ glow discharge, the average pressure in the presence of a beam of 100 mA is estimated to be that shown in Fig. III.38. Here the average pressure in one lattice cell is plotted as a function of the beam energy for three cases:

a) bending magnets and quadrupoles are fitted with distributed ion pumps and discharge-cleaned to give $\eta \approx 10^{-6}$ everywhere;

b) to reduce the aperture of the quadrupoles the integrated ion pump has been left out, resulting in a reduced effective pumping speed and in an increased desorption yield assumed to be $10^{-5}$, i.e. after several months of operation; and

c) $\eta$ of at least $10^{-7}$ after bakeout and a short conditioning period with the beam.

It can be seen that only the first method gives the required beam lifetime.

5.4.3 Some problems related to synchrotron radiation

The critical energy of the synchrotron radiation in CHEEP ranges from 0.574 keV at 5 GeV injection energy to 48.8 keV at 25 GeV. At this top energy -- based on the PETRA vacuum chamber profile -- about 25% of the total synchrotron radiation power of 5.2 MW will escape from the vacuum chamber by Compton scattering. A large fraction of radiated power would
be absorbed in the iron yokes of the magnets and thus contribute to the heating of the tunnel. Since the maximum admissible power which can be taken by the existing tunnel ventilation system is limited to 360 kW, the synchrotron radiation has to be absorbed by a sufficiently thick lead screen in good thermal contact with the water-cooled vacuum chamber.

The required thickness of the lead screen has to be determined not only for the absorption of the synchrotron-radiation power, but also to reduce to an acceptable level the integrated radiation dose in the magnet coils, and furthermore to limit to tolerable levels the production of toxic gases such as ozone or nitrogen dioxide.

Because of the closed iron yoke and the proposed construction of the H-type magnets the radiation escaping to the outside can be neglected, and toxic gases produced in the tunnel air are several orders of magnitude below standard tolerance levels of 0.1 ppm for ozone and 5 ppm for NO₂. The Compton-scattered radiation and the necessary radiation screen to absorb the radiation have been estimated in a similar way to that for EPIC⁴³. Assuming 3 mm of lead, the fraction of synchrotron power which is not removed by the cooling system of the vacuum chamber amounts to less than 10 kW. The integrated dose to the magnet coils for a 10-year period remains under these conditions at an acceptable level of 10⁵ rad. Figure III.39 shows the production rate of ozone in ppm per m inside the air space of the dipole magnet. This calculation is based on a G-value (number of molecules produced per 100 eV energy deposition) of 6 for ozone⁴⁴. Because the lifetime of ozone molecules is less than one hour, the expected equilibrium concentration should remain below 10 ppm up to electron energies of 25 GeV. The dangers of nitrogen dioxide which, together with humidity in the air, will form nitric acid, require further study. The initial production rate of this acid is determined by a G-value between 1.5 and 3, and hence at least a factor of 2 lower than for ozone.

Nitric acid corrosion of the iron yoke or other metal parts inside the magnets has to be evaluated in more detail, in particular with respect to long-term effects. Since the corrosion problems do not differ from those in PETRA or PEP because of the very similar synchrotron radiation characteristics and vacuum chamber designs, equivalent protective measures should be taken in CHEEP as in the other storage rings. Because the vacuum chamber section has been chosen slightly larger than is required for beam considerations, an additional lead screen could be introduced inside the magnet gap, thus reducing the amount of out-scattered radiation still further.

5.4.4 Special vacuum problems encountered in CHEEP

Apart from the vacuum system in the normal cells, several special cases have to be considered:

i) Polarization regions

The synchrotron radiation distribution and the maximum power density on the vacuum chamber produced by the strong magnetic field of the kinks has been estimated as 1.5 kW/cm² and seems to be within the range of what can be tolerated with the vacuum chamber adopted². However, adequate pumping of these areas poses severe problems because in some modes of operation the field in the bending magnets adjacent to the kinks is insufficient for ion pumping, even with the large cell diameter of the distributed ion pumps. A possible solution would be to mount a large number of independent
small ion pumps along the vacuum chamber at typical intervals of about 40 cm; nine such pumps per magnet are needed to produce the equivalent effective pumping speed of the distributed pumps in the normal bending magnets. For these regions it is proposed to use C-type magnets which would permit easy access to the vacuum chamber for pump con-
nexion. If the normal H-type were to be used, the magnet yoke would have to be pierced with holes for these connections.

ii) Interaction regions
Within the interaction region the vacuum will be improved with respect to the magnet
cells by installing a vacuum chamber with larger conductance, permitting the use of
large pumps. In this region the vacuum chamber will not be exposed to synchrotron
radiation, and only thermal desorption will determine the base pressure, expected to
be in the low $10^{-11}$ Torr range.

On either side of the intersection point a succession of vertical and horizontal bends
will steer the electron beam from its normal orbit to the level of the proton beam. The
vacuum chamber in these magnets will have to be considerably enlarged to accept the
beam, and consequently the conductance will be large so that again lumped pumps can
be used between the successive magnets. The design of this part of the vacuum system
should permit baking at temperatures of up to 300°C. Owing to the vertical bending,
synchrotron radiation will no longer be directed towards the cooled outer wall of the
vacuum chamber, and hence special cooling channels will have to be provided where needed.

6. THE PROTON BEAM

6.1 Introduction

The 60 proton bunches required in the SPS are provided in three PS cycles, each producing
a pulse of 20 bunches at 14 GeV/c. One pulse contains $0.7 \times 10^{13}$ particles, which is well
within the present performance of the PS. At the end of each PS cycle, the bunches are trans-
ferred to 20 SPS buckets in such a way that the SPS is filled with 60 equidistant bunches
after the third PS cycle. The total injection time, being equal to two PS cycles, is about
2.0 sec $^{45}$. No change in PS operation up to 14 GeV/c is required, although the bunch area
should stay preferably below 15 mrad in units of $\Delta \phi \cdot \text{RF-rad}$, which is reasonable at an in-
tensity of $0.7 \times 10^{13}$ particles/pulse.

The protons are accelerated in the SPS up to the extraction energy, at a somewhat re-
duced rate owing to the high beam intensity; they can collide continuously with the electrons
from 145 GeV upwards. Below this energy the protons have a revolution frequency differing
too much from that of the electrons for the bunches to be synchronized. At low energy, the
protons are separated from the electrons by means of a vertical bump to avoid excessive
proton tune shifts.

6.2 Transfer from the PS to the SPS

In order to diminish space-charge effects (see Section 7 of this chapter) in the SPS
the protons are transferred at 14 GeV/c, the highest momentum at which the existing transfer
channel TT10 can be operated. At this energy a short flat top is inserted in the PS cycle
and the bunches are transferred to the SPS in the sequence shown below (bunch to bucket
transfer).
Table III.6

Sequence of transfer

1  3  5  7  9  11  13  15  17  19
12 14 16 18 20  2  4  6  8 10

Since the individual bunches, about 15 nsec long in the PS, have to fit into the 5 nsec long SPS buckets, they must be compressed prior to extraction. This is done by shaping them on the unstable fixed point of the stationary RF bucket and subsequent rotation around the stable fixed point as indicated in Fig. III.40. This method has been successfully tested in the PS by D. Boussard, who obtained bunches of 5 nsec length with the required intensity.

After compression, all 20 bunches must be transferred within a lapse of time that is short compared to a quarter of a phase-oscillation period, i.e. within less than 500 μsec, otherwise they will continue to rotate in the bucket and lengthen again. Given the sequence shown above, this requires a kicker with a very high repetition rate, the time between two kicks corresponding to 11 bunch-to-bunch distances in the PS, i.e. 1.12 μsec, leading to a total ejection time of 22 μsec, which meets the above requirements.

Obviously, the present fast kicker (FAK) of the PS is not suitable because its pulse generator has too low a repetition rate. In order to overcome this limitation, J.-C. Schnuriger proposed a pulse generator which has some similarity to the one used for continuous ejection from the PS. It consists of a series of pulse-forming networks (PFN) each followed by a switch as indicated in Fig. III.41. The switches would be closed one after the other starting with the downstream one, producing in this way a train of current pulses, the pulse width being determined by the pulse-forming network length and the pulse interval by the switch triggering. It is obvious that the switches must remain conducting until the last pulse has passed, without deteriorating the shape of the pulses; it is felt that such a performance is best obtained from spark gaps which have a relatively long de-ionization time compared to thyratrons. Furthermore, they will not fire when the discharge of the neighbouringPFN makes the full voltage suddenly appear across the switch, a feature which cannot be obtained from multistage thyratrons. Finally, the rise-time of a spark gap compares favourably with that of a thyratron. Any loss in pulse-height by the residual resistance of the switches can be taken into account by the proper adjustment of the appropriate charging voltage. A small-scale model of such a pulse generator comprising two spark gaps and two PFNs has been built and tested by J.-C. Schnuriger. Delaying the second pulse up to 100 μsec did not produce any appreciable deterioration of its rise- and fall-times, which were around 25 nsec. A model consisting of five PFNs plus switches operating at 80 kV is being built.

According to P. Lefèvre, a kick strength of 300 G·m is sufficient provided that a septum, 3 mm thick, is used in conjunction with the usual kick enhancement. Thus two modules, identical to the FAK modules, operated with a PFN voltage around 80 kV, will provide the necessary deflection. Since it is virtually impossible to connect a second pulse generator to FAK and to retain its full availability for other users of the PS, it is felt that a new kicker consisting of nine modules should be provided for CHEEP operation. In order to save pulse generators, the modules will be grouped in pairs, each pair connected in series via a cable
providing a delay of 1.12 μsec, as indicated in Fig. III.41. Four groups are foreseen to limit the number of spark gaps per group. A ninth module is desirable as a spare. The kicker is operated according to the scheme outlined in Table III.7 for the first six bunches.

### Table III.7

<table>
<thead>
<tr>
<th>Bunch No.</th>
<th>Pulse source</th>
<th>Module active</th>
<th>Pulse source</th>
<th>Module active</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Group</td>
<td>PFN</td>
<td></td>
<td>Group</td>
</tr>
<tr>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>1 module of FAK</td>
</tr>
<tr>
<td>12</td>
<td>2</td>
<td>1</td>
<td>3</td>
<td>1</td>
</tr>
<tr>
<td>5</td>
<td>3</td>
<td>1</td>
<td>5</td>
<td>2</td>
</tr>
<tr>
<td>14</td>
<td>4</td>
<td>1</td>
<td>7</td>
<td>3</td>
</tr>
<tr>
<td>5</td>
<td>1</td>
<td>2</td>
<td>1</td>
<td>4</td>
</tr>
<tr>
<td>16</td>
<td>2</td>
<td>2</td>
<td>3</td>
<td>1</td>
</tr>
<tr>
<td>etc.</td>
<td>etc.</td>
<td>etc.</td>
<td>etc.</td>
<td>etc.</td>
</tr>
</tbody>
</table>

The rise- and fall-times of the kick pulse must each fall within 90 nsec, which is well within the performance of FAK modules. Once the PFNs further upstream in a group have to be used, which may cause some deterioration of the rise-times, alternate bunches will already have been ejected.

The fast kicker in the SPS clearly requires the same high repetition rate and the same considerations as above. The new kicker magnet developed in the SPS for multibatch transfer has the required number of modules. Its rise-time is less than 110 nsec \(^3\), well within the required 360 nsec. Since we cannot accommodate a second kicker magnet in the SPS for CHEEP operation, pulse generators must be changed when switching from normal SPS operation to CHEEP operation. It is proposed to arrange the ten modules of the magnet in five groups, the two modules of each group being fed in parallel from one pulse generator consisting of a series of 20 PFNs and spark gaps (or grounded grid thyatron). It is clear that this gives a sizeable number of elements. Detailed study and tests are needed before assessing not only the performance but also the reliability and lifetime of such a system delivering nearly \(10^7\) shots per year. Table III.8 summarizes the parameters for the fast transfer.

There are two possible options which would avoid high pulse repetition rates of the kickers in the PS and SPS, and at the same time would reduce the space-charge tune shift at injection in the SPS.

One option would be to upgrade the equipment required for transfer and injection into the SPS to a momentum of 26 GeV/c. At this momentum, the PS RF system can hold a bunch area of up to 12 mrad in a bunch which is not longer than the SPS bucket, and the tune shifts would be less than 0.08 at injection in the SPS.

The other option is to distribute the PS bunch over several, at least two, 200 MHz buckets prior to ejection in order to reduce the line density. The existing 200 MHz RF system
of the PS can handle those sub-bunches because the voltage requirement is reduced by the splitting up. The tune shift is a function of the number of sub-bunches; however, already in the worst case, i.e. only two sub-bunches at 14 GeV/c, the tune shift would not exceed 0.12.

Both options have in common that the transfer would not need to be fast. In the PS, the present FAK, after minor control modifications, could deliver a bunch to the SPS every 6 msec in the sequence indicated in Table III.6. On the SPS side, a new pulse generator would nevertheless be needed; this would be a relatively simple device consisting of five PFNs plus five thyatrons only, and a resonant charging power supply.

Unfortunately, both options also have in common that they require substantial new hardware or hardware modifications. What is required in the first case is rather obvious and will not be discussed further. In the second case, the sub-bunches have to be recombined in the SPS before interaction with the electrons. This can be done by an auxiliary low-frequency RF system as described below.

Since small bunches of comparable phase-space density are less stable than bigger ones, care must be taken to ensure the longitudinal stability of the sub-bunches. This can best be done by splitting the PS bunch into only two sub-bunches.

After "slow transfer" and acceleration to 150 GeV/c, RF voltage will be reduced until the bunches touch each other in the stationary buckets. Pulsing of the auxiliary RF system, operating at a low frequency, will rotate the bunches in synchrotron phase space until one sits on top of the other. The 200 MHz RF, which must be switched off and compensated during
the rotation, will then be switched on again with its maximum voltage and both bunches will
be trapped in one bucket. Figure III.42 shows rotation and trapping of one of the two sub-
bunches.

Filamentation after recapture will blow up the bunch area. This is welcome because
longitudinal stability is improved, although the blow-up must stay within limits imposed by
the acceptance of the RF system and by the acceleration rate required after compression.
Since the acceptance increases at higher energy it is proposed to perform bunch compression
at the highest possible energy, i.e. just before collisions with the electrons start. Thus
a high acceleration rate is possible before compression; after compression, it has to be
reduced somewhat, but the time is not wasted as e-p collisions already take place.

The parameters of such an auxiliary RF system are listed in Table III.9. The relatively
high frequency and the pulsed operation for only about 3 msec will ease its implementation.
The total cavity length is around 5 to 6 m and the diameter would be the same as the diameter
of the 200 MHz cavities.

| Table III.9 |
| Parameter of auxiliary RF system |

| Harmonic number | \( h = 1320 \) |
| Frequency | \( f = 57.25 \text{ MHz} \) |
| Voltage for bunch rotation | \( U = 2 \text{ MV} \) |
| Time for rotation (small amplitude) | \( t_0 = 2.65 \text{ msec} \) |
| Actual time of rotation | \( t = 1.11 \times t_0 = 2.94 \text{ msec} \) |
| Blow-up factor | \( \frac{A_{\text{bucket}}}{A_{\text{bunch}}} = 2 \) |

Figure III.42 is based on these parameters. The final bucket has an area of \( 2.8 \Delta \phi \text{·RF-rad} \).

6.3 Beam handling in the SPS

Since the SPS improvement programme will be implemented by the time e-p operation is
expected to start, we assume that the RF system consists of four cavities, each having four
sections; each cavity is fed from two transmitters of 500 kW each\(^{15}\).

The RF voltage at 200 MHz required for trapping after fast transfer must create a sta-
tionary bucket with a height exceeding the momentum spread of the bunch. The minimum voltage
must then be about 4.5 MV, and the bunch area will be diluted by some small factor which helps
to stabilize the bunch. Since the RF system can provide a maximum voltage of 8.6 MV at 14 GeV/c,
the beam can easily be trapped, leaving a sufficient margin for beam loading and acceleration.

In fact, an acceleration rate of 58 GeV/sec can be obtained immediately after injection
with a phase angle of 9° in the limit of vanishing beam loading. Reduction of the bucket
area by space-charge forces turns out to be negligible. During the front porch the phase angle
can be increased rapidly; for example, the maximum possible acceleration rate at 20 GeV/c
already exceeds the standard rate of 167 GeV/sec. Thus the voltage limits the rate only very
close to injection. At higher energy, two more limitations have to be examined: the rate of
energy transfer from the transmitters to the beam via the cavities\textsuperscript{[7]}, and the longitudinal high-frequency instability\textsuperscript{[8]}. The instability seems to require a continuously growing bucket area, thus limiting the stable phase angle for a given voltage.

The energy which is supplied to the cavities by the transmitters between the passage of two bunches is

\[ W = \frac{P_{tr}}{(60 \ f_0)} = 1.5 \ J, \]

where \( P_{tr} = 4 \ \text{MW} \) is the power of the transmitters. In order to take into account the power fed back from the beam, we assume, following the experience gained from present SPS operation\textsuperscript{[9]}, that only one quarter of this energy (0.38 J) is accepted by the passing bunch. This energy has to cover the energy gain per turn as well as possible losses which occur through excitation by the bunched beam of higher-order parasitic modes in the cavities and in the rest of the machine. By scaling from the measurements in SPEAR\textsuperscript{[9]}, an estimate of these losses can be obtained in a way similar to that used for PEP. This indicates that these losses are very small (\( 3.8 \times 10^{-6} \ J/\text{turn} \)) since the bunch is relatively long and the peak current is low. Neglecting the parasitic losses, we find that the acceleration rate could rise to 300 GeV/sec before being limited by this effect, leaving a useful safety margin.

The second limitation turns out to be the more serious one if our assumptions are correct. As explained in more detail in Section 7, the bunch, driven by the instability, will most likely increase in area and provide enough local energy spread to stabilize itself. If we assume that the beam loading does not substantially decrease the accelerating voltage, then the stable phase angle must be limited to \( \approx 25^\circ \) (\( \Gamma = 0.4 \)) in order to provide a bucket large enough to contain the growing bunch. This leads to a maximum acceleration rate of 90 GeV/sec. Table III.10 summarizes the RF parameters.

<table>
<thead>
<tr>
<th>Table III.10</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>RF parameters</strong></td>
</tr>
<tr>
<td>------------------</td>
</tr>
<tr>
<td>Frequency</td>
</tr>
<tr>
<td>Number of cavities</td>
</tr>
<tr>
<td>Sections per cavity</td>
</tr>
<tr>
<td>Transmitters per cavity</td>
</tr>
<tr>
<td>Nominal power of one transmitter</td>
</tr>
<tr>
<td>Filling time of cavity</td>
</tr>
<tr>
<td>Maximum RF voltage at 14 GeV/c</td>
</tr>
<tr>
<td>RF voltage required for trapping</td>
</tr>
<tr>
<td>Initial stable phase angle</td>
</tr>
<tr>
<td>Effective RF voltage during acceleration</td>
</tr>
<tr>
<td>Stable phase angle during acceleration</td>
</tr>
<tr>
<td>Bunch area at 400 GeV/c</td>
</tr>
<tr>
<td>Momentum spread at 400 GeV/c ( \sigma_p/E )</td>
</tr>
<tr>
<td>Transition parameter ( \gamma_{tr} )</td>
</tr>
</tbody>
</table>
Figure III.4 shows an example of an SPS cycle based on this acceleration rate. After waiting for two PS cycles the beam is accelerated during 4.4 sec to 400 GeV where a flat top of = 2 sec is inserted for slow and fast ejection. The power dissipation averaged over the cycle is less than the maximum tolerable one (36 MW). This cycle would provide an average of $2 \times 10^{12}$ protons/sec. nearly 60% of that provided by a standard cycle in the future.

It is obvious that one can invent a wide variety of cycles, e.g. some which include a flat part for doing e-π physics at constant energy. However, it is felt that this is beyond the scope of this report because the at present unknown future requirements of experimentation with fixed targets will strongly influence the shape of the cycle.

As mentioned before, the two beams will be separated during acceleration to avoid excessive proton tune shifts at low energy. This is accomplished by means of a vertical bump in the proton orbit, which is generated by two correction dipoles separated by half a betatron wavelength and disposed symmetrically around the interaction point. The field in the dipoles will be roughly constant over the cycle. Thus the bump amplitude will decrease with increasing proton energy, and the bump will have collapsed at higher energy as required. The amplitude of the bump is obtained from the conditions that the protons should never experience a tune shift higher than $\Delta Q_p = 0.01$. For a Gaussian density distribution in the electron beam, the proton tune shift is reduced by a factor $f$ if the protons are displaced vertically by

$$\Delta z = \sqrt{2\pi} \sigma_{ze}^*,$$
valid if
$$\Delta z / \sigma_{ze}^* \geq 3.$$

Since
$$\Delta Q_p \sim E_p^{-1},$$
the factor $f$ is given by
$$f = E_{p,\text{min}} / E_{p,\text{inj}},$$

$E_{p,\text{min}}$ is the lowest proton energy in the cycle, where the beams can be brought together, and $E_{p,\text{inj}}$ is the proton energy at injection. Clearly, the maximum displacement is required for $E_{p,\text{min}} = 400$ GeV yielding $\Delta z = 3$ mm. In order to have a sufficient safety margin, a further displacement corresponding to $3 \sigma_{zp}^*$ at 14 GeV/c should be added yielding finally $\Delta z = 7.5$ mm for the maximum bump amplitude.

Fast extraction of the beam does not appear to present any particular problems. The momentum spread of the bunches $\pm 0.8 \times 10^{-3}$ is within the momentum bite accepted by the extraction system, i.e. $\pm 1 \times 10^{-3}$.

Slow extraction appears to be more difficult because the strongly bunched beam must be debunched within a reasonable time. Then a careful excitation by resonant kicks must be applied to bring the particles close to the electrostatic system where they are peeled off from the beam. The debunching time equals approximately the time $T_d$ it takes for the coasting bunches to close the gap between them:

$$T_d = \left( \frac{\Delta p}{p} f_3 \right)^{-1} = 0.17 \text{ sec}.$$
It could be argued that longitudinal instabilities will lengthen this time. However, in the ISR, where it is established that the growth of the bunch area is dominated by the instabilities, no anomaly is observed during debunching, and the final momentum spread corresponds to the momentum spread of the beam in the RF bucket. It is clear that only experiments in the SPS will reveal whether this comparison is justified.

The resonance mechanism which is used to provide the desired beam growth is not affected by the presence of the insertion. However, the insertion has not the full acceptance required for the large-amplitude particles which have received the final kick immediately before being extracted. Table III.11 gives the typical betatron amplitudes, normalized to $\beta_x = 106 \text{ m}$, at the electrostatic septum during the last turns for the three resonance modes used at present.

<table>
<thead>
<tr>
<th>Order of resonance</th>
<th>Third Quasi-integer</th>
<th>Half-integer</th>
<th>Integer</th>
</tr>
</thead>
<tbody>
<tr>
<td>$Q_x$</td>
<td>$26 \frac{2}{3}$</td>
<td>26.55</td>
<td>26.9</td>
</tr>
</tbody>
</table>

Amplitude in mm at turn:

\[
\begin{align*}
\text{at turn:} & \\
\left\{ \begin{array}{c}
n - 2 \\
n - 1 \\
\end{array} \right. & \\
\{ \begin{array}{c}
34 \\
40 \\
45 \\
\end{array} & & 32 \\
\{ \begin{array}{c}
39 \\
50 \\
\end{array} & & 38 \\
\end{align*}
\]

Although a detailed study of the extraction has not yet been made, it seems possible from inspection of the present layout that the elements generating the steps in amplitude can be placed downstream from the insertion and before the electrostatic septum, at least for the slow extraction towards the North Hall. Examination of Table III.11 indicates that in this case the amplitude would not exceed 40 mm (at $\beta_x = 106$), reducing the required acceptance to $E_x = 15\pi \times 10^{-6} \text{ rad\cdotm}$ which is the nominal acceptance of the insertion. Figure III.43 shows the aperture ($E_x = 15\pi \times 10^{-6} \text{ rad\cdotm}$) required by the proton beam in the region of the insertion septa during extraction. The other beam, being bent away by common magnets and septa, is either an electron or positron beam at 25 GeV. The aperture and position of the septa are compatible with the configurations of the proton beam at an injection momentum as low as 14 GeV/c; this is shown in Figs. III.10 and III.11. A closed-orbit distortion, scaled from 5 mm peak-to-peak at $\beta_x = 106 \text{ m}$, is added, yielding the over-all envelope shown as full line. It is clear that the insertion as proposed here means a substantial reduction of the SPS acceptance from the standard $E_x = 50\pi \times 10^{-6} \text{ rad\cdotm}$ to $15\pi \times 10^{-6} \text{ rad\cdotm}$ during e-p operation, a fact which would have a serious impact on the SPS flexibility for fixed-target operation.

If we do not operate in e-p mode we can always retract the septa into the aperture of the electron ring, and retrace the insertion as discussed in Section 3.1 of this chapter. Figure III.44 shows the proton beam envelope for $E_x = 50\pi \times 10^{-6} \text{ rad\cdotm}$ passing through the rearranged septa. Then the limitations are eliminated and the SPS should function as if the insertion were not there.
7. BEAM STABILITY

7.1 Proton beam

7.1.1 Introduction

It is assumed that the proton bunches have a parabolic line density; this seems to be a natural bunch form. Some estimates about the broad-band impedance in the SPS have been made based on scaling from measurements in the ISR\(^{58,51}\) and on beam observations in the SPS\(^{52}\). The longitudinal high-frequency impedance \(Z_L\) divided by the mode number \(n = f_{\text{inst.}}/f_0\) is estimated to be

\[
|Z_L/n| \approx 2\pi \Omega
\]

with an inductive and resistive part of about equal magnitude (14 \(\Omega\)). At very low frequencies the impedance is determined by the skin effect and can be larger.

For a round chamber the transverse impedance is related to the longitudinal one by

\[
Z_T = \frac{2\pi}{2\pi \rho b^2 n} Z_L,
\]

with \(b\) being the chamber radius. Approximating the SPS vacuum chamber by a round pipe of radius \(b = 0.022\) m, we get for the vertical transverse impedance at high frequencies \(|Z_T| = 90\) MHz/m. There may also be resonators in the ring which present large impedances with narrow bandwidth, but they are not yet well known.

7.1.2 Microwave instabilities (bunch lengthening)

This instability shows up as a blow-up of the momentum spread and bunch length and is connected with the appearance of microwave signals. It can be explained as a coasting beam instability inside bunches, occurring if the frequency of the instability is high enough that the corresponding wavelength is short compared to the bunch length, and if the growth rate (calculated for the coasting beam) is large compared to the phase oscillation frequency. The coasting beam stability criterion can be applied using the instantaneous momentum spread and current in the bunch\(^{59,61}\)

\[
\left( \frac{Z_L}{n} \right) \leq \frac{3F m e^2 c^2 |n| (\Delta\gamma)^2}{\epsilon Y}.
\]

For bunches with a parabolic line density the quantity \((\Delta\gamma)^2/I\) has the same value everywhere along the bunch, and the form factor is about \(F = 0.65\). The bunch is more stable if its area is large. For stationary conditions (\(\Gamma = 0\)) at 270 GeV/c we find a minimum bunch area \(A_b = 1.2 \Delta\gamma\cdot\text{RF-rad}\) necessary to get stability in the presence of a longitudinal impedance \(|Z_L/n| = 20\) \(\Omega\) and for an RF voltage of \(U = 5\) MV. This gives a bunch of full length \(\ell = 0.72\) m or \(\ell/c = 2.4\) nsec. During acceleration \(\gamma\), \(\gamma\) changes; however, it can be shown that we always have stability for a full bucket with \(\Gamma = 0.4\) and \(U = 5\) MV. This bunch (or bucket) area is shown in Fig. III.45 as a function of energy. Around transition energy a Q-jump will be necessary to avoid instability. The instability will probably always adjust the bunch area without much overshoot. At injection a somewhat larger bucket area (\(\Gamma < 0.4\)) is necessary to allow for trapping and matching of the newly injected beam.
7.1.3 **Longitudinal coupled bunch mode instabilities**

The phase oscillations executed by different bunches can be coupled together by an impedance with a memory, usually a resonator. Such a coupled bunch mode can have an exponentially growing amplitude. For $k_b$ equidistant bunches with total current $I_s$, a resonator with shunt resistance $R_s$, resonant frequency $\omega_r$, and quality factor $Q'$ will cause such an instability having an approximate maximum growth rate $1/\tau$ \cite{59}:

$$
1/\tau \approx \frac{\omega_s R_s I_s f_{\omega} D_1}{\cos \phi_s R_s f_0} \frac{1}{2\pi (f_r/c)^2 + 2},
$$

\[(III.18)\]

where $\omega_s$ is the phase oscillation frequency. The factor $D_1$ depends on the attenuation of the induced resonator signal between the passage of adjacent bunches; it is

$$
D_1 \approx 2ae^{-\alpha} \quad \text{for} \quad \alpha = \frac{\pi f_r}{Q'k_b f_0} \gg 1
$$

and

$$
D_1 \approx 1 \quad \text{for} \quad \alpha \ll 1.
$$

The spread $S$ in phase oscillation frequencies $\omega_s$ can provide Landau damping and stabilize the coupled bunch oscillations. We can give an approximate stability criterion,

$$
S \geq 4|1/\tau| + |\Delta\omega_s|.
$$

The second term $\Delta\omega_{st}$ is the shift of the incoherent frequencies due to the inductive wall impedance $Z_{L,\text{ind}}$ \cite{59}:

$$
\Delta\omega_{st} = \frac{3\omega_s I_s (2\pi R/c)^3}{2\pi^2 k_b h \cos \phi_s} \frac{Z_{L,\text{ind}}}{n},
$$

\[(III.19)\]

which is $\Delta\omega_{st} \approx 240 \text{ sec}^{-1}$ for our assumed impedance. The frequency spread $S$ due to the non-linearity of the standard RF wave form is

$$
S = \omega_s \phi_b/16 \approx 185 \text{ sec}^{-1}
$$

\[(III.20)\]

with $\phi_b$ being half of the full bunch length measured in terms of RF phase angle. This spread is not quite large enough to provide stability; however, the SPS will have installed a higher harmonic cavity, increasing the spread $S$ considerably. Assuming twice the spread (III.20), a growth rate $1/\tau \approx 20 \text{ sec}^{-1}$ as calculated from (III.18) could still be stabilized. A narrow bandwidth impedance of up to $R_s = 2 \Omega$ could therefore be tolerated. Most of the realistic resonators will have a broad bandwidth and are therefore less dangerous because the induced signal will decay considerably between the passage of adjacent bunches. The bandwidth of the RF cavity includes frequencies which can excite coupled bunch modes. If this is a problem, a feedback system could cure it relatively easily since only a limited number of low frequencies are involved.

7.1.4 **Single-bunch head-tail instability**

This instability is caused by a short-range interaction between particles in one bunch and does not need a wall impedance with a long memory. Particles at the head of the bunch
executing betatron oscillations can excite the particles at the tail. Half a phase oscillation later, the latter particles are at the head and excite the particles now at the tail. For this to produce a growing (or damped) oscillation it is necessary that the phases of the betatron oscillations change while the particles move from the tail to the head or vice versa. Such a phase change is provided if the chromaticity

\[ \xi = \frac{\Delta Q}{Q} \frac{\partial p}{\partial p} \]

is different from zero. The lowest head-tail mode is damped for \( \xi > 0 \) above transition energy and for \( \xi < 0 \) below transition energy. The stability criterion is more complicated for the higher modes because they are governed by the frequency dependence of the transverse impedance\(^{54}\). For a relatively flat impedance the instability of these higher modes should be very weak; it has actually not been observed in most machines. For \( \xi = 0 \) all modes are stable; however, it is difficult to keep the chromaticity exactly zero at all energies. Furthermore, a non-zero chromaticity could be beneficial to stabilize the transverse coupled bunch mode instability.

For a smooth transverse resistive impedance \( Z_{TR}(\omega) \) we can give the growth (or damping) rate of the lowest head-tail mode for non-zero chromaticity:

\[ \frac{1}{\tau_1} = \text{sign} \left( \frac{d f_\xi}{d p} \xi \right) \frac{e^{\xi c^2 Z_{TR}(\omega)}}{4\pi f_0 \mu_0 c^2 \gamma k_B^2}, \]

with \( \omega_\xi = 2\pi f_0 Q \xi / |\eta| \).

Above transition energy \( d f_\xi / d p \) is negative and we get damping for a positive chromaticity.

7.1.6 Transverse coupled bunch mode instability

If the transverse impedance has a memory the head-tail oscillations executed by different bunches can be coupled together. In this case the lowest head-tail mode is no longer stable for \( \xi = 0 \) but has a growth rate\(^{54}\)

\[ \frac{1}{\tau_2} = \frac{e^{\xi c^2 F'(\xi) Z_{TR}(\omega)}}{4\pi f_0 \mu_0 c^2 \gamma 2 n R}, \]

with \( \omega_n = 2\pi(n + Q)f_0, F' \) being a form factor which is 0.8 for \( \xi = 0 \) and smaller for larger \( \xi \). To get the total growth rate we have to add the single bunch part (III.21) to the coupled mode part (III.22). Since the first part can be made negative (damping) by choosing a non-zero chromaticity of the correct sign, we can compensate the coupled mode contribution and get over-all stability for the lowest mode. The higher modes, however, could be slightly unstable and may have to be damped with octupoles.

7.1.8 Incoherent Q-shift

The incoherent Q-shift is given by\(^{55}\)

\[ \Delta Q = \frac{-(2\pi R)^2 r_q k_F^*}{\frac{p}{(\pi E_{\gamma} / \gamma)}(1 + \frac{k_F}{E_{\gamma}}) \gamma^2 k_B^2}. \]

Here \( r_q \) is the classical proton radius and \( \gamma E_{\gamma} \) is the normalized transverse emittance. The factor \( F^* \) takes into account the effect of the image fields: it is \( \approx 1.2 \) between 14 and 20 GeV.
If bunch compression is performed in the PS, the tune shift and, unfortunately, also the tune spread, become 0.25 in the SPS at injection (14 GeV/c). Such a tune shift occurs at injection during normal operation of the PS booster and is well tolerated by the beam without emittance growth. Tune shifts of 0.3 and more can be handled also in the PS at injection after careful compensation of resonances\(^{56}\). Admittedly, a big ring like the SPS might be more sensitive.

If bunch compression can be performed in the SPS at 150 GeV, as was discussed in Section 6 of this chapter, the maximum tune shift at injection will be reduced to 0.12 or less. The number still seems fairly high in the light of present experience with injection at 10 GeV/c. However, it is known that the SPS tolerates tune variations of this magnitude at 200 GeV/c, and we may therefore hope that a similar tolerance may be achieved at 14 GeV/c, especially after compensation of the most dangerous resonances.

7.1.7 Intrabeam scattering

The growth times due to intrabeam scattering\(^{57}\) at injection, where the effect is strongest, are listed below

\[
\tau_E = 6 \times 10^2 \text{ sec} \\
\tau_X = 4 \times 10^3 \text{ sec} \\
\tau_Z = 2 \times 10^4 \text{ sec}
\]

It is obvious that this effect is negligible.

7.2 Electron beam

7.2.1 Introduction

For the electron ring the impedance will be lower than for the SPS. We assume a high-frequency impedance about four times smaller than that of SPEAR but with the same frequency dependence. For the inductive part we assume \(Z_{L,\text{ind}} = 10\). The following parameters are used (Table III.12).

<table>
<thead>
<tr>
<th>E (GeV)</th>
<th>(U_0) (MeV/turn)</th>
<th>(\alpha_E) (sec(^{-1}))</th>
<th>(\sigma_{\text{D/E}})</th>
<th>(U_{\text{RF}}) (MV)</th>
<th>(Q_S)</th>
<th>(\sigma_{\text{SO}}) (cm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>0.38</td>
<td>3.3</td>
<td>(10^{-3})</td>
<td>10</td>
<td>0.06</td>
<td>5.7</td>
</tr>
<tr>
<td>25</td>
<td>4.66</td>
<td>81</td>
<td>(8 \times 10^{-4})</td>
<td>67</td>
<td>0.06</td>
<td>2.4</td>
</tr>
</tbody>
</table>

At high energies these parameters are given by the lattice alone, whilst at low energy, they are determined by the lattice and wiggler magnets. The RF voltage of 10 MV at low energy has been chosen to allow for bunch lengthening.
7.2.2 Bunch Lengthening

The increase in bunch length and momentum spread has been estimated using a scaling law \( \sigma_s = 5.3 \times 10^{-6} \xi^{1/3.2} \), based on SPEAR data and corrected for the expected lower impedance in CHEEP:

\[
\xi = \frac{I_s |n| R^3}{k_B Q_s^2 E}
\]

(I\(_s\) in mA, R in m, and E in GeV).

We get the following results:

<table>
<thead>
<tr>
<th>E (GeV)</th>
<th>( \sigma_s ) (cm)</th>
<th>( \sigma_s / \sigma_{so} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>10.8</td>
<td>1.9</td>
</tr>
<tr>
<td>25</td>
<td>3.2</td>
<td>1.3</td>
</tr>
</tbody>
</table>

This bunch lengthening is tolerable and should not affect the performance.

7.2.3 Longitudinal coupled bunch mode instabilities

To calculate these instabilities we approximate the Gaussian bunch form of the electrons by a parabolic longitudinal line density with the full length

\[
\xi = 2\sqrt{2} \sigma_s
\]

For the short electron bunch at 25 GeV the frequency shift due to the inductive wall impedance (III.19) is probably large enough to eliminate Landau damping, and we are left with radiation damping only. From Eq. (III.18) we find that in this case a narrow bandwidth impedance with \( R_s = 8 \text{ M}\) is still tolerable. At 5 GeV the bunch is longer, we get a smaller frequency shift due to the inductive impedance and a larger spread in phase oscillation frequencies. The Landau damping is quite strong and an impedance of about 18 M\( \Omega \) can be tolerated.

7.2.4 Transverse instabilities

The smaller impedance and the radiation damping make this instability less dangerous for electrons. Based on Eqs. (III.21) and (III.22) we can estimate that the transverse instability should occur only in the presence of very high impedance resonators.

8. ELECTRON-POSITRON INJECTOR

It is proposed to fill the CHEEP electron/positron storage ring (e-ring) in the SPS tunnel from a fast-cycling synchrotron which in turn is fed from a linac.

A reasonable injection energy for the e-ring is 5 GeV, which is about the top energy of NINA, the electron synchrotron in Daresbury Laboratory, UK. Since NINA was phased out recently and since its magnets may become available, it was natural to consider a synchrotron making use of NINA's combined function magnets. Typical problems of the injection system could then be studied thus helping to establish realistic performance figures.
It is not claimed that this is an optimum solution. However, at the moment it is not apparent either that a specially designed synchrotron would offer decisive advantages as far as performance is concerned.

A possible site in CERN for the injector is shown in Fig. III.2. The electrons are brought to the SPS via existing transfer tunnels; only the short bend between TT2a and TT60 might require some civil engineering work.

8.1 Injector Synchrotron

8.1.1 Output Beam Emittance

Injection into the e-ring will be in the radial phase plane, which sets an upper limit for the emittance of the synchrotron output beam, since repeated injection into a given bucket is necessary. Fast kicker magnets in the e-ring will bring the stored beam close to a thin septum inflector magnet which separates it from the incoming beam. To ensure efficient injection the incoming beam emittance should be less than 25% of the acceptance available in the main ring for radial betatron oscillations. This sets \(E(2\sigma) \leq 3.75\pi \times 10^{-6} \text{ rad m}\).

In a combined-function electron synchrotron with F and D magnets of equal length, the radial betatron oscillations are antidamped\(^{59}\) and the guide field must be fast cycled to minimize the radial beam emittance. For a given cycling rate the radial emittance will increase with increasing output energy and, in the unmodified NINA lattice using 50 Hz excitation to 5 GeV, the radial emittance is a factor of 3 to 4 greater than the above limit. Several methods are available for modifying the natural radial damping\(^{59}\), which may be characterized by the value of the radial damping partition number \(J_x\). Figure III.46 shows the theoretical radial betatron emittance \(E_x(2\sigma)\) as a function of \(J_x\) and peak energy in the NINA magnet lattice. The magnet excitation is of the form \(B(t) = B_0 + B_2(1 - \cos \omega t), \omega = 100\pi \text{ sec}^{-1}\), and the initial electron (positron) emittance corresponds to the emittance expected at the linac output, as explained in Section 8.2 below. It may be seen that by operating at a peak energy somewhat below 5 GeV, together with a modification of \(J_x\) from its natural value of \(-0.92\) to \(J_x \geq 0\), the constraint \(E_x \leq 3.75\pi \times 10^{-6} \text{ rad m}\) may be satisfied.

8.1.2 Lattice

The NINA synchrotron consisted of 20 alternating-gradient combined-function FODO cells in a ring of mean radius 35.1 m\(^{60}\). For the injector synchrotron, the long and short straight sections are increased from 3.5 m to 3.9 m, and 1.0 m to \(\sim 1.09\) m, respectively, to give a mean radius of 36.67 m which is 1/30 of the mean radius of the e-ring. The four pairs of programmable quadrupoles are powered to give nominal betatron tunes of \(Q_x = 5.71, Q_z = 5.76\); and the required modification of \(J_x\) is achieved by the inclusion, in each of four 3.9 m straights, of combined-function triplet magnets. These are D-profile magnets with the outer elements of length 0.375 m bending in the same sense as the main guide field, whilst the centre element, of length 0.75 m, bends in the reverse sense. To obtain \(J_x = 0\) the magnet pole profile must provide \((1/B)(dB/dx) = 6 \text{ m}^{-1}\). The RF frequency is chosen to be identical with that of the main ring at 312.5 MHz. The four RF cavity units, each consisting of three \(3/2\) cells, are placed in symmetrical 3.9 m straights and are powered by a single klystron.

Injection into the booster is from an electron (positron) linac as described in Section 8.2, the inflector magnet being situated in one of the eight remaining long straight sections. At
peak energy, the beam is extracted using a fast kicker and a septum magnet, which are positioned in adjacent long straight sections separated in betatron phase by \( \approx 90^\circ \).

A sketch of the lattice geometry and betatron functions is presented in Fig. III.47 and the main parameters are listed in Table III.13 below.

<table>
<thead>
<tr>
<th>Parameters of injector synchrotron</th>
</tr>
</thead>
<tbody>
<tr>
<td>Repetition frequency</td>
</tr>
<tr>
<td>Peak energy</td>
</tr>
<tr>
<td>Mean radius</td>
</tr>
<tr>
<td>Magnet bending radius</td>
</tr>
<tr>
<td>Number of superperiods</td>
</tr>
<tr>
<td>Betatron tunes ( Q_x, Q_z )</td>
</tr>
<tr>
<td>Momentum compaction factor</td>
</tr>
<tr>
<td>RF harmonic number</td>
</tr>
<tr>
<td>RF frequency</td>
</tr>
<tr>
<td>RF voltage at peak energy (( \phi_s = 45^\circ ))</td>
</tr>
<tr>
<td>RF structure shunt impedance</td>
</tr>
<tr>
<td>Total RF structure length</td>
</tr>
<tr>
<td>Peak RF structure power</td>
</tr>
<tr>
<td>Synchrotron tune ( Q_s ), at peak energy</td>
</tr>
<tr>
<td>Revolution time</td>
</tr>
<tr>
<td>Radial damping partition number, ( J_x )</td>
</tr>
<tr>
<td>Positron radial emittance at peak energy ( E_{x/n(2\sigma)} )</td>
</tr>
<tr>
<td>Momentum spread at peak energy (( \pm 2\sigma ))</td>
</tr>
</tbody>
</table>

8.2 Linac injectors, intensities, and filling times

8.2.1 Electron filling

The circulating current required in the e-ring is 104 mA \( \approx 1.5 \times 10^{13} \) particles contained in a total of 60 equidistant bunches. To provide this current we use an S-band electron linac capable of accelerating a peak current of 4 A to 100 MeV. At this intensity level, the linac output momentum spread will be \( \approx 1.5\% \), which exceeds the momentum acceptance of the synchrotron; so, for electron filling, we accelerate only 1 A to perhaps 110-120 MeV and estimate the emittance \( E(2\sigma) \leq 2\pi \times 10^{-6} \) rad\cdot m and momentum spread \( \leq \pm 5 \times 10^{-3} \).

The longitudinal acceptance is limited by the peak momentum spread, which occurs within the first millisecond after injection. If this is limited to 1\% then the corresponding allowable phase extent of the injected bunch is approximately 140\(^\circ\) of the synchrotron RF cycle. In order to match the linac pulse to the acceptance of a bucket in the synchrotron, the gun is pulse-modulated to produce 1.4 nsec long pulses. At an average current of 1 A in such a pulse, which appears as a bunch train at the linac output, one traps \( 9 \times 10^9 \) electrons per synchrotron bucket. The number of buckets one can fill in the synchrotron is determined by
the booster repetition frequency and the transverse damping time in the e-ring. The latter sets a lower limit to the time which must elapse between two successive injections into the same bucket of the e-ring. In order to speed up the filling, the damping time in the main ring is shortened from 3.4 sec to 0.6 sec by the kink magnets together with a 37 m long set of wiggler magnets\(^\text{11}\) operating at 0.6 T. Thus two bunches can be accelerated in the synchrotron, and the 60 buckets in the e-ring will be full after 22 sec if an injection efficiency of 75% is assumed.

### 8.2.2 Positron filling

For positron production, the linac output beam is directed on to a tungsten target, the full 4 A capability being used. The conversion factor \(e^+/e^-\) is estimated to be \(\approx 1.8 \times 10^{-3}\)\(^\text{61}\). The resulting positrons are accelerated by a further linac section to an energy of 200 MeV, at which energy the beam will be contained within a transverse emittance of \(\approx 23 \times 10^{-6}\) rad·m with a momentum spread of \(\pm 1.5\%\).

This emittance is within the acceptance of the synchrotron, if we assume the same vacuum chamber dimensions as in NINA; however, the momentum spread is a factor of 3 too big. To overcome this, an energy compression system\(^\text{62}\) is used, which consists essentially of a system of magnets to produce a path length proportional to the particle momentum followed by a suitably phased RF correction cavity to reduce the momentum spread. The power for this cavity may be tapped off the linac transmitter. A theoretical improvement of a factor of 15-20 in momentum spread is possible; however, we assume that only a factor of 5 will be achieved, yielding a momentum spread of \(\pm 3 \times 10^{-3}\).

With this initial condition and the injection energy higher for positrons than for electrons, the maximum allowable phase extent is \(\approx 240^\circ\) of the booster RF cycle \(\equiv 6\) linac RF cycles. The linac electron gun is thus pulse-modulated to produce a 4 A pulse of 2.2 nsec duration which, with the conversion factor above, gives the expected positron intensity as \(10^8\) e\(^+\) per booster bunch, a factor of 100 lower than with electrons. Hence the filling time becomes 40 min, which still looks very reasonable when compared with the expected beam lifetime of about 20 hours.

After the initial phase of running, which probably involves a great deal of electron-proton experiments, it may become desirable to shorten the positron filling time. Several options can be offered to accomplish this, some of them at moderate cost, but all cumulative in effect:

i) increase of the gun current to 20 A,

ii) use of a subharmonic prebuncher,

iii) increase of primary electron energy,

iv) positron accumulator ring after linac,

v) second RF system in synchrotron,

vi) bunch accretion in the e-ring.

i) Gun currents up to 20 A are used in S-band linacs of low duty cycle\(^\text{61}\). Such high current would reduce the e\(^+\) filling time by a factor of 5.

ii) High gun currents may be easier to handle if combined with a subharmonic prebuncher.

Such a device would be placed between gun and linac, and would compress a long gun pulse to the required length\(^\text{64}\), increasing the charge contained in the accepted pulse by about a factor of 2.
iii) Increase of the energy of the primary electrons would give a proportional gain in positron yield.\(^{63}\)

iv) The insertion of a small positron accumulator between the linac and the synchrotron, like the PIA ring in DESY, would be another possibility.\(^{65}\) A similar ring was proposed to reduce the filling time in LEP.\(^{66}\) However, such a device would demand a substantial increase of the positron linac energy in order to function with a sufficiently short damping time.

v) A longer linac bunch train would be trapped, if the acceptance of the synchrotron were increased. This could be achieved by using a low-frequency RF system at low energy. Assuming a frequency of 52 MHz (500 kV) for this system, it is found that a linac bunch train more than three times longer can be trapped, reducing the filling time by the same factor. After acceleration to about 1 GeV, when the bunch area is "reduced" by adiabatic damping, the 312 MHz system takes over. A fifth harmonic is superimposed on the magnet cycle of the synchrotron to provide a reduced \(\hat{B}\) at the take-over point.

iv) If we filled 240 or 300 equidistant buckets in the e-ring instead of 60 in the same time, and combined them into 60 by means of a by-pass, a factor of 4 or 5 could be gained in filling time at the expense of an extraction system at 5 GeV and a new 200 m long tunnel. The synchrotron would accelerate 8 or 10 bunches instead of 2. After the 240 or 300 buckets in the e-ring were filled, bunch accretion would be performed by ejecting appropriate single bunches into the by-pass, where they would be delayed by 3 to 4 bunch distances, and by combining them with the bunches in the 60 master buckets. Recombination would take less than 4 sec. A possible layout of such a by-pass is shown in Fig. III.48. Existing tunnels and transfer lines TT10 and TT60 could be used; the only addition would be a 200 m long linking tunnel and transfer line.

9. LAYOUT, UTILITIES, AND BUILDINGS

9.1 Layout

The shape of the underground ring tunnel housing the SPS is shown in Fig. III.1. It consists of six circular arcs connected by long straight sections (LSS). In the circular sextants the tunnel has an inside diameter of 4 m, while in the LSS it widens to 6 m and 8.5 m. For the SPS, LSS 1 contains the injection system, LSS 2 and LSS 6 contain the extraction systems to the North and to the West Experimental Areas, respectively. The RF accelerating structure is in LSS 3 and the internal dump in LSS 4. The only relatively free LSS is LSS 5, which is therefore considered for the interaction region between the SPS proton and the electrons/positrons. For the e-ring, LSS 6 is to be used for the injection system and LSS 2 for the RF cavities required for operation at 25 GeV. For operation at 30 GeV, a second RF system will be required, which could be placed in LSS 6.

As has already been shown, the e-ring and the SPS lattices are identical. The e-ring elements -- bending magnets, quadrupoles, and correcting lenses -- will have the same azimuthal positions as the corresponding SPS elements. The electron beam will be in a plane situated 110 cm above the SPS. This distance was chosen for ease of handling and installation of the various machine components, and for sufficient independence of the two machines and of their services. The offset of 110 cm also satisfies the requirement of the geometry in the interaction region.
Figure III.49 shows the situation in the 4 m diameter sextants of the tunnel. Several solutions have been shown to be possible for the support system of the e-ring elements. The solution shown in this figure consists of suspending the machine from bolts anchored in the roof of the tunnel. Several tests have shown this solution to be feasible and to provide a sufficient safety factor.

Also shown in Fig. III.49, next to the e-ring magnet are the supplementary cable trays, the cooling pipes, and the water-cooled cables required to power the magnets.

The situation in the LSS is much more complicated because of the large amount of specialized equipment of the SPS. The LSS also connect to the pits through which all the services (pipes and cables) are brought into the tunnel from the surface Auxiliary Buildings. It has not yet been possible to carry out any detailed design for these areas. It is, however, envisaged that the e-ring beam, while remaining at the same height, will be offset radially by about 54 cm, so as to avoid the SPS equipment. The most difficult situation is in LSS 3, where the SPS RF system, with its auxiliary equipment, occupies a very large part of the tunnel.

The only LSS for which a layout has been taken to a first design stage is LSS 5, the interaction region. The situation is shown in Fig. III.50, where it can be seen that, with the geometry considered at present, the machine only just fits into the tunnel. It will be necessary to recalculate a slightly modified geometry and to design the magnetic elements to very tight dimensions before tackling the problems of supports and services in this area.

9.2 Utilities

9.2.1 Electrical power supply

The total power requirements for the e-ring amount to about 30 MVA for operation at 25 GeV. This figure takes into account the efficiencies of the various systems, covering all the losses in transformer power supplies, etc.

The CERN Prevesin site is supplied with electrical power by a 380 kV overhead line, connecting at Geneva to the French super-grid. This line, capable of transmitting power in excess of 500 MVA, feeds the main electrical substation through three 90 MVA transformers, supplying the 18 kV distribution system of the SPS complex*

Two of these transformers are connected to networks feeding all the pulsed loads of the SPS and of the North Experimental Area. The third transformer is used to supply the general services and all the other steady loads on the site. The loading on this transformer is very small compared to its rating of 90 MVA, which has been chosen to allow its use as a spare in case of breakdown of any of the other transformers. The e-ring, being essentially a d.c. machine, can be connected to the steady-load network, the corresponding transformer being able to carry the full load.

Depending on the detailed distribution of the load, it might become necessary to strengthen the cables leading from some of the Auxiliary Buildings to the electrical substation. It will also be necessary to install in each Auxiliary Building supplementary MV/LV transformers and the related busbar and switching equipment.

*) The installations described here will exist, and will be in operation by the end of 1978.
9.2.2 Cooling system

The new machine will require two cooling systems: one for the magnets, the other to evacuate the energy deposited in the walls of the vacuum chamber by synchrotron radiation.

The cooling system for the SPS magnets presents sufficient reserve to provide the cooling of the e-ring magnets.

The detailed implementation is easy because of the identity of the lattices of the two machines. The magnet system of the electron machine can be connected directly to the same inlet/outlet branching-off points on the main demineralized-water pipes used for the SPS magnets.

The existing cooling plants and distribution system are, however, unable to cope with the heat generated by synchrotron radiation. It will therefore be necessary to make provision for a new cooling system.

The heat evacuated from all the machine elements -- SPS and e-ring -- is transferred via heat exchangers to the primary, raw water which is brought to the site from the lake of Geneva and fed by a pumping station to the ring mains serving the SPS. A study is being carried out to check whether this raw-water system can handle the supplementary load from CHEEP and to analyse various methods of increasing its capability.

9.3 Buildings

All the services, power-supplies, control equipment, etc., for the SPS are housed in the six Auxiliary Buildings placed at the top of the access pits. While no detailed layout study has yet been carried out, it is obvious that these buildings cannot provide the necessary space required for CHEEP.

All six buildings will have to be enlarged by about 200 m\(^2\) for MV/LV switchgear, cooling plant, power supplies, racks for electronics, computer and peripherals, etc.

Moreover, a supplementary Auxiliary Building will be required next to Auxiliary Building No. 2 for the installation of the RF plant. The area for the RF plant is tentatively estimated to require about 2000 m\(^2\).

10. CONTROL SYSTEM

The control system of the SPS has been shown to be efficient, flexible, and reliable. The control system for CHEEP will therefore make use, whenever possible, of the same hardware modules and the same software. This will involve the installation of a new set of computers linked, via CAMAC and a multiplex system, to the e-ring equipment. This system must initially be independent of the SPS, since it must be possible to continue the operation of the SPS during the construction and commissioning periods of the e-ring. The concept of multicomputers interacting via a message-switching system makes it possible to interconnect eventually the two systems, so that the SPS and the e-ring can eventually be controlled from one common control centre.

11. SURVEY

As in any accelerator, the survey of CHEEP will be done in two successive steps, the first one being the provisional installation of the accelerator components and the second one the final positioning of these elements. These two operations must be done as quickly as possible and particularly during any SPS shutdowns, using the SPS metrology.\(^67\).
The only prerequisite is the possibility to re-measure the long geodetic chain of braced-quadrilaterals by means of the reference marks on the brackets, which were installed on the inner and the outer walls of the tunnel during the construction of the SPS. These measurements must be carried out before the provisional installation of the first magnet.

The installation will start using the brackets on the walls implying that the plane at the level of the brackets and the reference marks on the quadrupoles must be completely free of any obstacle during installation, and after in order to be able to make changes to the SPS components, if necessary. It is possible that there may be obstacles in a few areas, e.g. at the radio-frequency tanks, but such complications must be avoided as much as possible to shorten the installation time and not to lose the accuracy reached by the normal geometric configuration.

The first operation will be the layout of the supports and the provisional positioning of the quadrupoles and dipoles. Such a solution is feasible because of the identity of the lattice of the two machines. It is possible from the survey point of view to accept changes in the dipole size, but the reference marks fixed to the quadrupoles of the e-ring should be on the vertical above those fixed on the SPS quadrupoles. The supports of the quadrupoles must be designed to allow for a maximum accuracy to be achieved in a minimum of time during the final, micrometric adjustments.

It is too early at this stage of the study to freeze the design of the supports. But it must be pointed out that the supports for the quadrupoles must be completely decoupled from those of the dipoles. Using the tridimensional referential system of the SPS, the final installation of the quadrupoles must be carried out first, and that of the dipoles only where at least three quadrupoles are already in place.

After the final survey, it is necessary to plan the smoothing of the local curvature for successive sets of three adjacent quadrupoles of CHEEP. The operation of smoothing is very important but very short in time. The relative radial error of the quadrupoles will be the same as in the SPS, e.g. 0.10 mm. For the levelling, the relative accuracy of 0.10 mm attained for one quadrupole relative to the two adjacent quadrupoles is a reasonable expectation.

It should be appreciated that all the numerous detection and correction elements must be aligned to the same precision. This is a time-consuming operation and requires a careful design to facilitate adjustment.

12. PARAMETER LISTS

12.1 Performance

<table>
<thead>
<tr>
<th>Parameter</th>
<th>p-ring</th>
<th>e-ring a)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Luminosity (peak)</td>
<td>$0.5 \times 10^{12}$</td>
<td>cm$^{-2}$ sec$^{-1}$</td>
</tr>
<tr>
<td>Nominal energy</td>
<td>270</td>
<td>25</td>
</tr>
<tr>
<td>Total number of particles</td>
<td>$2 \times 10^{13}$</td>
<td>GeV</td>
</tr>
<tr>
<td>Circulating current</td>
<td>140</td>
<td>$1.5 \times 10^{13}$</td>
</tr>
<tr>
<td>Number of bunches</td>
<td>60</td>
<td>104</td>
</tr>
<tr>
<td>Circumference</td>
<td>$2\pi \times 1100$</td>
<td>mA</td>
</tr>
<tr>
<td></td>
<td>$2\pi \times 1100.013$</td>
<td>m</td>
</tr>
</tbody>
</table>

a) Operating with four orbit kinks
### Performance (Contd.)

<table>
<thead>
<tr>
<th>Parameter</th>
<th>p-ring</th>
<th>e-ring</th>
<th>rad·m</th>
<th>mm</th>
</tr>
</thead>
<tbody>
<tr>
<td>Emittance (b) E_x/\pi</td>
<td>(20 \times 10^{-6}/\beta_Y)</td>
<td>(1.1 \times 10^{-7})</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Emittance (b) E_z/\mu</td>
<td>(10 \times 10^{-6}/\beta_Y)</td>
<td>(0.77 \times 10^{-8})</td>
<td></td>
<td></td>
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<tr>
<td>Beam size at (\sigma_x)</td>
<td>0.34</td>
<td>0.39</td>
<td>mm</td>
<td></td>
</tr>
<tr>
<td>intersections (\sigma_z)</td>
<td>0.072</td>
<td>0.059</td>
<td>mm</td>
<td></td>
</tr>
<tr>
<td>Bunch length (\sigma_s)</td>
<td>300</td>
<td>30</td>
<td>mm</td>
<td></td>
</tr>
<tr>
<td>Energy spread (\sigma_T/E)</td>
<td>(0.8 \times 10^{-3})</td>
<td>(1.4 \times 10^{-3})</td>
<td></td>
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</tr>
<tr>
<td>Beam-beam (\Delta q_X)</td>
<td>0.008</td>
<td>0.006</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Tune shifts (\Delta q_z)</td>
<td>0.008</td>
<td>0.014</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Energy range for collisions</td>
<td>150 - 400</td>
<td>5 - 25(30)</td>
<td>GeV</td>
<td></td>
</tr>
<tr>
<td>Lifetime</td>
<td>SPS cycle</td>
<td>20 h</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Filling time ((p e^-/e^+)</td>
<td>1.8 sec</td>
<td>0.5/40 min</td>
<td></td>
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</tr>
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#### 12.2 Insertion

<table>
<thead>
<tr>
<th>Parameter</th>
<th>p-ring</th>
<th>e-ring</th>
</tr>
</thead>
<tbody>
<tr>
<td>Amplitude (f_x) functions</td>
<td>6.5</td>
<td>1.5</td>
</tr>
<tr>
<td>(f_z) functions</td>
<td>0.6</td>
<td>0.3</td>
</tr>
<tr>
<td>Dispersion (D_x) functions</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td>(D_z) functions</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td>Crossing angle (2\phi)</td>
<td>(\pm 2.5)</td>
<td>mrad</td>
</tr>
<tr>
<td>Free space</td>
<td>(\pm 5)</td>
<td>m</td>
</tr>
</tbody>
</table>

#### 12.3 Electron ring

##### 12.3.1 Lattice

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
<th>Unit</th>
</tr>
</thead>
<tbody>
<tr>
<td>Circumference</td>
<td>6911.586</td>
<td>m</td>
</tr>
<tr>
<td>Phase advance/cell</td>
<td>90°</td>
<td>m</td>
</tr>
<tr>
<td>Amplitude functions (\beta_x) max/min</td>
<td>105/19</td>
<td>m</td>
</tr>
<tr>
<td>(\beta_z) max/min</td>
<td>105/19</td>
<td>m</td>
</tr>
<tr>
<td>Dispersion function (D_x) max</td>
<td>3</td>
<td>m</td>
</tr>
<tr>
<td>Bending radius</td>
<td>741.3</td>
<td>m</td>
</tr>
<tr>
<td>Length of one cell</td>
<td>63.996</td>
<td>m</td>
</tr>
<tr>
<td>Nominal tune</td>
<td>28</td>
<td>m</td>
</tr>
<tr>
<td>Momentum compaction</td>
<td>(1.66 \times 10^{-3})</td>
<td></td>
</tr>
<tr>
<td>Uncorrected chromaticity (\varepsilon_x/\varepsilon_z)</td>
<td>-36/-36</td>
<td>m</td>
</tr>
<tr>
<td>Variation of damping (\partial \beta_x/(\partial E/E))</td>
<td>-135</td>
<td>m</td>
</tr>
<tr>
<td>Total number of cells</td>
<td>108 - 6</td>
<td>(one insertion)</td>
</tr>
<tr>
<td>Cells with kink magnets</td>
<td>16</td>
<td>m</td>
</tr>
<tr>
<td>Cells with wiggler magnets</td>
<td>5</td>
<td>m</td>
</tr>
<tr>
<td>Closed-orbit distortion at (\beta_{\text{max}})</td>
<td>(\pm 2.5 \times 10^{-3})</td>
<td>m</td>
</tr>
</tbody>
</table>

Length and disposition of elements \(\equiv\) SPS; electron orbit 1.1 m above SPS orbit. The missing dipoles are moved by one half-cell away from the LSS.

---

*a) Operating with four orbit kinks; b) Protons: \(E/\pi = (2\sigma)^2/\beta;\) Electrons: \(E/\pi = \sigma^2/\beta;\) throughout parameter list.
### 12.3.2 Beam characteristics derived from lattice

**Operation at**

- a) 25 GeV
- b) 25 GeV with four kinks
- c) 30 GeV

<table>
<thead>
<tr>
<th>Parameter</th>
<th>a</th>
<th>b</th>
<th>c</th>
</tr>
</thead>
<tbody>
<tr>
<td>Synchrotron energy loss/turn</td>
<td>46.6</td>
<td>51.8</td>
<td>96.7</td>
</tr>
<tr>
<td>Max. emittance $\sigma_{x_0}/\beta_x$</td>
<td>$1.0 \times 10^{-7}$</td>
<td>$1.1 \times 10^{-7}$</td>
<td>$1.4 \times 10^{-7}$</td>
</tr>
<tr>
<td>Natural bunch length $\sigma_s$</td>
<td>24</td>
<td>31</td>
<td>23</td>
</tr>
<tr>
<td>Natural energy spread $\sigma_E/E$</td>
<td>$0.8 \times 10^{-3}$</td>
<td>$1.4 \times 10^{-3}$</td>
<td>$0.9 \times 10^{-3}$</td>
</tr>
<tr>
<td>Damping time $\tau_x$</td>
<td>25</td>
<td>22</td>
<td>14</td>
</tr>
<tr>
<td>Polarization time $\tau_p$</td>
<td>100</td>
<td>30</td>
<td>40</td>
</tr>
<tr>
<td>Max. long. polarization</td>
<td>76</td>
<td>76</td>
<td>?</td>
</tr>
</tbody>
</table>

### 12.3.3 RF parameters

a, b, and c same meaning as before

<table>
<thead>
<tr>
<th>Parameter</th>
<th>a</th>
<th>b</th>
<th>c</th>
</tr>
</thead>
<tbody>
<tr>
<td>Orbital frequency</td>
<td>43.38</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Radiofrequency</td>
<td>312.31</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Harmonic number</td>
<td>7200</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Active length of cavities</td>
<td>96</td>
<td>96</td>
<td>192</td>
</tr>
<tr>
<td>Cavity shunt impedance</td>
<td>25</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Group velocity/c</td>
<td>≈ 0.1</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Electron/positron current</td>
<td>0.1</td>
<td></td>
<td>A</td>
</tr>
<tr>
<td>RF peak voltage</td>
<td>69</td>
<td>74.8</td>
<td>138</td>
</tr>
<tr>
<td>Overvoltage factor</td>
<td>1.48</td>
<td>1.44</td>
<td>1.43</td>
</tr>
<tr>
<td>Quantum lifetime</td>
<td>100</td>
<td></td>
<td>h</td>
</tr>
<tr>
<td>Stable phase angle</td>
<td>45°</td>
<td>44°</td>
<td>45°</td>
</tr>
<tr>
<td>Synchrotron tune $Q_s$</td>
<td>0.062</td>
<td>0.064</td>
<td>0.079</td>
</tr>
<tr>
<td>Synchrotron frequency</td>
<td>2.7</td>
<td>2.8</td>
<td>3.4</td>
</tr>
<tr>
<td>Loss by synchrotron radiation by fundamental mode of cavity</td>
<td>4.9</td>
<td>5.4</td>
<td>10.1</td>
</tr>
<tr>
<td>Total RF generator power</td>
<td>8</td>
<td>8-9</td>
<td>16</td>
</tr>
<tr>
<td>Efficiency of RF generator</td>
<td>60</td>
<td></td>
<td>%</td>
</tr>
<tr>
<td>Number of cavities</td>
<td>16</td>
<td>16</td>
<td>32</td>
</tr>
<tr>
<td>Number of active cells/cavity</td>
<td>12</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Phase advance/cavity cell</td>
<td>$\pi/2$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Physical length/cavity</td>
<td>110</td>
<td>110</td>
<td>220</td>
</tr>
<tr>
<td>Inside diameter of cavity</td>
<td>$&gt; 750$</td>
<td></td>
<td>mm</td>
</tr>
<tr>
<td>Beam hole $\varnothing$</td>
<td>$&gt; 100$</td>
<td></td>
<td>mm</td>
</tr>
</tbody>
</table>
12.3.4 Vacuum system in lattice

Vacuum chamber

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Conductance</td>
<td>67 ml/sec</td>
</tr>
<tr>
<td>Dimensions</td>
<td>114 x 56 mm^2</td>
</tr>
<tr>
<td>Bake-out temperature</td>
<td>150 °C</td>
</tr>
<tr>
<td>Thermal out-gassing</td>
<td>8 x 10^-9 Torr</td>
</tr>
<tr>
<td>Molecular desorption yield</td>
<td>10^-6 molecules/phot-o-e</td>
</tr>
<tr>
<td>Pumping speed of linear pump</td>
<td>55 1./(sec.m)</td>
</tr>
<tr>
<td>Pumping speed of lumped pump</td>
<td>40 1/sec</td>
</tr>
</tbody>
</table>

Synchrotron radiation 25(30) GeV

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Critical energy</td>
<td>47(81) keV</td>
</tr>
<tr>
<td>Linear power density</td>
<td>1.1(1.9) kW/m</td>
</tr>
<tr>
<td>Maximum power density on wall</td>
<td>160(280) W/cm^2</td>
</tr>
<tr>
<td>Gas desorption from wall</td>
<td>0.58(0.66) x 10^-6 Torr</td>
</tr>
</tbody>
</table>

Bremsstrahlung on residual gas

| Lifetime at 25 GeV        | 27 h           |
| Lifetime at 5 GeV         | 20 h           |

The vacuum chamber is made from aluminium; it is in situ glow-discharge-cleaned everywhere.

Linear pumps are installed in bending magnets and quadrupoles. The system in the lattice is designed for operation at 30 GeV with 0.1 A.

12.3.5 Magnet system

Parameters of standard elements at 30 GeV

<table>
<thead>
<tr>
<th>Quadrupole</th>
<th>Dipole (Version H)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Over-all length</td>
<td>3210 mm</td>
</tr>
<tr>
<td>Over-all width</td>
<td>400 mm</td>
</tr>
<tr>
<td>Over-all height</td>
<td>400 mm</td>
</tr>
<tr>
<td>Core weight</td>
<td>1300 kg</td>
</tr>
<tr>
<td>Coil weight</td>
<td>300 kg</td>
</tr>
<tr>
<td>Total weight</td>
<td>1700 kg</td>
</tr>
<tr>
<td>Magnetic length</td>
<td>3000 mm</td>
</tr>
<tr>
<td>Bore Ø/Gap height</td>
<td>159.5 mm</td>
</tr>
<tr>
<td>Field at pole</td>
<td>0.1228 T</td>
</tr>
<tr>
<td>Gradient</td>
<td>1.54 T/m</td>
</tr>
<tr>
<td>Number of turns/pole</td>
<td>3</td>
</tr>
<tr>
<td>Excitation current</td>
<td>1306 A</td>
</tr>
<tr>
<td>Size of conductor</td>
<td>25 x 16 mm^2</td>
</tr>
<tr>
<td>Inductance/magnet</td>
<td>8.76 x 10^-4 H</td>
</tr>
<tr>
<td>Cooling channels/conductor</td>
<td>1</td>
</tr>
<tr>
<td>Hole diameter</td>
<td>4 mm</td>
</tr>
</tbody>
</table>
Parameters of standard elements at 30 GeV (Contd.)

<table>
<thead>
<tr>
<th></th>
<th>Quadrupole</th>
<th>dipole (Version H)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Water pressure drop</td>
<td>5</td>
<td>5</td>
</tr>
<tr>
<td>Water flow rate</td>
<td>6</td>
<td>4.3</td>
</tr>
<tr>
<td>Temperature rise</td>
<td>15</td>
<td>15</td>
</tr>
<tr>
<td>Total number</td>
<td>216 - 11</td>
<td>744 - 15</td>
</tr>
<tr>
<td>Total resistance</td>
<td>0.77</td>
<td>1.89</td>
</tr>
<tr>
<td>Total power dissipation</td>
<td>1.51</td>
<td>3</td>
</tr>
</tbody>
</table>

The elements are designed for operation at 30 GeV.

Parameters of special magnets in insertion

<table>
<thead>
<tr>
<th>Magnet function</th>
<th>Number</th>
<th>Yeke type</th>
<th>Max. field (T)</th>
<th>Length (m)</th>
<th>Vertical aperture (mm)</th>
<th>Horizontal aperture (mm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Common separating bend</td>
<td>4</td>
<td>S</td>
<td>0.165</td>
<td>1.9</td>
<td>130</td>
<td>450/150</td>
</tr>
<tr>
<td>Double-sided septum</td>
<td>4</td>
<td>S</td>
<td>0.823</td>
<td>1.9</td>
<td>130</td>
<td>490/250</td>
</tr>
<tr>
<td>Low-beta quadrupole</td>
<td>8</td>
<td>ELB</td>
<td>10.1 T/m</td>
<td>1.0</td>
<td>130/85</td>
<td>90/190</td>
</tr>
<tr>
<td>Polarization bends</td>
<td>8</td>
<td>EH</td>
<td>0.3</td>
<td>2.8</td>
<td>190</td>
<td>250</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Orbit restoring vertical</td>
<td>16</td>
<td>EV</td>
<td>0.36</td>
<td>2.8</td>
<td>190</td>
<td>250</td>
</tr>
<tr>
<td>Orbit restoring large aperture</td>
<td>5</td>
<td>EH</td>
<td>0.22</td>
<td>6.26</td>
<td>190</td>
<td>250</td>
</tr>
<tr>
<td>Orbit restoring double standard</td>
<td>3</td>
<td>EL</td>
<td>0.22</td>
<td>6.26</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Parameters of quadrupoles in insertion

<table>
<thead>
<tr>
<th>Type</th>
<th>Magnetic length (m)</th>
<th>Field gradient (T/m)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Q₀</td>
<td>1.0</td>
<td>-10.1</td>
</tr>
<tr>
<td>Q₁</td>
<td>1.0</td>
<td>9.15</td>
</tr>
<tr>
<td>Q₂</td>
<td>1.0</td>
<td>-1.53</td>
</tr>
<tr>
<td>Q₃</td>
<td>1.0</td>
<td>-2.73</td>
</tr>
<tr>
<td>Q₄</td>
<td>1.0</td>
<td>1.50</td>
</tr>
<tr>
<td>Q₅</td>
<td>3.045</td>
<td>-2.81</td>
</tr>
<tr>
<td>Q₆</td>
<td>3.045</td>
<td>2.06</td>
</tr>
<tr>
<td>Q₇</td>
<td>3.045</td>
<td>-1.23</td>
</tr>
<tr>
<td>Q₈</td>
<td>3.045</td>
<td>1.47</td>
</tr>
</tbody>
</table>
Parameters of kink magnets

Length 0.4 m
Max. field 2.1 T
Aperture standard
Total number 16 (4 + 12)

Parameters of injection wiggler magnets

<table>
<thead>
<tr>
<th>Type</th>
<th>Short S</th>
<th>Long L</th>
</tr>
</thead>
<tbody>
<tr>
<td>Length/element</td>
<td>1 m</td>
<td>2 m</td>
</tr>
<tr>
<td>Elements/cell</td>
<td>2</td>
<td>3</td>
</tr>
<tr>
<td>Total number of elements</td>
<td>10</td>
<td>15</td>
</tr>
<tr>
<td>Arrangement/cell</td>
<td>S L L L S</td>
<td></td>
</tr>
<tr>
<td>Field</td>
<td>0.6 T</td>
<td></td>
</tr>
<tr>
<td>Bending radius</td>
<td>25.5 m</td>
<td></td>
</tr>
<tr>
<td>Aperture</td>
<td>standard</td>
<td></td>
</tr>
</tbody>
</table>

12.3.6 Injection system

Fast kicker + septum combined with fast bump

Injection energy 4.8 GeV
Filling time $e^-/e^+$ 0.5/40 min
Required damping time $\tau_X$ 0.6 sec
Number of kinks active all
Field strength in kink magnets 0.7 T
Kick pulse, length at bottom $< 370$ (65) a) nsec
length at top $> 5$ nsec
Injection efficiency 75 %

12.4 Electron injector

12.4.1 Linac

<table>
<thead>
<tr>
<th></th>
<th>$e^+$</th>
<th>$e^-$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Frequency</td>
<td>S-band</td>
<td></td>
</tr>
<tr>
<td>Energy</td>
<td>200</td>
<td>120</td>
</tr>
<tr>
<td>Conversion efficiency</td>
<td>$1.8 \times 10^{-3}$</td>
<td>-</td>
</tr>
<tr>
<td>Gun current averaged over 1 linac cycle</td>
<td>4</td>
<td>1</td>
</tr>
<tr>
<td>Linac cycles during gun pulse</td>
<td>6</td>
<td>4</td>
</tr>
<tr>
<td>Energy spread</td>
<td>$\pm 3 \times 10^{-3}$ b)</td>
<td>$\pm 3 \times 10^{-3}$</td>
</tr>
<tr>
<td>Emittance $E_x/\pi$</td>
<td>$5.75 \times 10^{-6}$</td>
<td>$0.5 \times 10^{-6}$</td>
</tr>
</tbody>
</table>

---

a) For bunch accretion
b) After energy compression
12.4.2 Synchrotron

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Maximum energy</td>
<td>5 GeV</td>
</tr>
<tr>
<td>Average radius</td>
<td>36.67 m</td>
</tr>
<tr>
<td>Nominal tune</td>
<td>5.75 m</td>
</tr>
<tr>
<td>Bending radius</td>
<td>20.77 Mm</td>
</tr>
<tr>
<td>Damping partition numbers</td>
<td>J_x = 0; J_z = 1; J_s = 3</td>
</tr>
<tr>
<td>Lattice</td>
<td>NINA + 4 damping triplets</td>
</tr>
<tr>
<td>RF</td>
<td>512.3 MHz</td>
</tr>
<tr>
<td>Harmonic number</td>
<td>240</td>
</tr>
<tr>
<td>Stable phase angle</td>
<td>45 degrees</td>
</tr>
<tr>
<td>RF peak voltage</td>
<td>3.5 MV</td>
</tr>
<tr>
<td>Power of RF generator</td>
<td>100 kW</td>
</tr>
<tr>
<td>Performance with</td>
<td>e^+ e^-</td>
</tr>
<tr>
<td>Number of particles/bunch</td>
<td>10^8 e 10^10</td>
</tr>
<tr>
<td>Number of bunches</td>
<td>2 e 2</td>
</tr>
<tr>
<td>Emittance E_x/\eta</td>
<td>1 x 10^-6 e 75 x 10^-6 rad m</td>
</tr>
<tr>
<td>Energy spread</td>
<td>\pm 1.4 x 10^-3 e \pm 1.4 x 10^-3</td>
</tr>
<tr>
<td>Ejection energy</td>
<td>4.8 e 4.8 GeV</td>
</tr>
<tr>
<td>Repetition rate</td>
<td>50 e 50 Hz</td>
</tr>
</tbody>
</table>

12.5 Requirements on SPS

12.5.1 New magnets

Parameters of special magnets in insertion

<table>
<thead>
<tr>
<th>Magnet</th>
<th>Length (m)</th>
<th>Gap height (mm)</th>
<th>Magnetic field p channel e (T)</th>
<th>Number</th>
</tr>
</thead>
<tbody>
<tr>
<td>Double septum S1 a)</td>
<td>1.9</td>
<td>140</td>
<td>0.825/-0.104 \pm 0.165</td>
<td>2</td>
</tr>
<tr>
<td>Double septum S2 a)</td>
<td>1.9</td>
<td>140</td>
<td>0.495 \pm 0.165</td>
<td>2</td>
</tr>
<tr>
<td>Common separating bend</td>
<td>1.9</td>
<td>102</td>
<td>\pm 0.165</td>
<td>4</td>
</tr>
</tbody>
</table>

a) These septa will also serve as the hadron calorimeter.

Parameters of quadrupoles in insertion

<table>
<thead>
<tr>
<th>Quadrupole</th>
<th>Magnetic length (m)</th>
<th>Gradient</th>
<th>Required aperture for beam</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>low-\beta mode</td>
<td>detuned mode</td>
</tr>
<tr>
<td></td>
<td></td>
<td>(T/m)</td>
<td>(T/m)</td>
</tr>
<tr>
<td>Q_1</td>
<td>5.5</td>
<td>-15.9</td>
<td>15.5</td>
</tr>
<tr>
<td>Q_2</td>
<td>5.5</td>
<td>-15.9</td>
<td>-16.7</td>
</tr>
<tr>
<td>Q_3</td>
<td>4.145</td>
<td>15.9</td>
<td>3.0</td>
</tr>
<tr>
<td>Q_4</td>
<td>4.145</td>
<td>15.9</td>
<td>-8.2</td>
</tr>
</tbody>
</table>
12.6.2 RF

Future SPS RF system: 4 cavities, 4 sections each
8 transmitters, 500 kW each

Frequency 200 MHz
Harmonic number 4620
Interaction length/cavity 16.45 m
Filling time of cavity 566 nsec
SPS long. impedance 2/n 20 Ω
Peak voltage 8.6 MV
Acceptance at 14 GeV a) 0.68 Δδγ•RF-rad
Bunch area at 14 GeV 0.42 Δδγ•RF-rad
Number of particles/bunch 3.3 × 10^{11}
Momentum spread at 400 GeV ±0.75 × 10^{-3}
Maximum acceleration 90 GeV/sec
Minimum acceleration time to 400 GeV 4.4 sec

12.6.3 Lattice (unchanged, cf. CERN document No. 1050 (1972)).

12.6.4 Injection

Mode of transfer Fast bunch to bucket transfer
Momentum 14 GeV/c
Number of kicker groups 5
Number of kicker modules/group 2 (in parallel)
Number of switches + PBN/group 20 (in series)
Repetition time of kick 1.12 μsec
Pulse length at top ≥ 10 nsec
Pulse length at bottom ≤ 360 nsec

12.6 Requirements on PS

12.6.1 General parameters

Number of PS pulses/SPS cycle 3
Number of bunches/pulse 20
Repetition time PS 0.9 sec

Beam characteristics at ejection
Momentum 14 GeV/c
Number of particles/pulse 0.7 × 10^{13}
Bunch area < 0.02 Δδγ•RF-rad
Bunch length (±20) ±2.5 nsec
Momentum spread (±20) ±3 × 10^{-3}
Max. RF voltage 200 kV

a) Vanishing beam loading.
12.6.2 Ejection

<table>
<thead>
<tr>
<th>Mode of transfer</th>
<th>Fast bunch to bucket transfer with a FAK</th>
</tr>
</thead>
<tbody>
<tr>
<td>Number of kicker groups</td>
<td>4</td>
</tr>
<tr>
<td>Number of modules/group</td>
<td>2 (in delayed series)</td>
</tr>
<tr>
<td>Number of spark gaps + PFN/group</td>
<td>5 (in series)</td>
</tr>
<tr>
<td>Number of spare modules</td>
<td>1</td>
</tr>
<tr>
<td>Total physical length</td>
<td>2452 mm</td>
</tr>
<tr>
<td>Strength of kick</td>
<td>0.03 T·m</td>
</tr>
<tr>
<td>Repetition time of kick</td>
<td>1.12 μsec</td>
</tr>
<tr>
<td>Pulse length at top</td>
<td>≥ 10 nsec</td>
</tr>
<tr>
<td>Pulse length at bottom</td>
<td>≤ 90 nsec</td>
</tr>
<tr>
<td>Duration of ejection</td>
<td>21.95 μsec</td>
</tr>
</tbody>
</table>

Acknowledgements

It is a pleasure to thank G. Saxon and T. Swain, who clarified our ideas on injection into the e-ring, and P. Lefèvre who explained to us the PS performance limits relevant to CHEEP. From helpful talks with H. Kuhn, J.C. Schnuriger, G. Schröker and W.C. Middelkoop we learnt more about fast kicker systems.

We also thank R. Billinge, D.A. Gray, J. Le Duff, K. Steffen, V.P. Suller and G.A. Voss for the stimulating discussions we had with them during the e-p study week at Rutherford Laboratory, whose hospitality we gratefully acknowledge.
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Fig. III.1  Schematic layout of SPS

Fig. III.2  Possible location of electron/positron injector
a) Luminosity versus crossing angle at nominal energy

b) Electron tune shifts versus crossing angle at nominal energy

Fig. III.3

Fig. III.4  Example of an SPS cycle during e-p operation
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a) Luminosity versus electron energy at nominal proton energy (270 GeV)

b) Relative electron current versus electron energy

Fig. III.6
a) Luminosity versus proton bunch length at nominal energy

b) Electron tune shifts versus proton bunch length at nominal energy

Fig. III.7

a) Luminosity at nominal energy versus coupling in e-ring

b) Proton tune shifts versus coupling in e-ring

Fig. III.8
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Fig. III.17 SPS proton ring structure functions - detuned low-beta, fully matched into lattice to give $Q_x = 26.58$, $Q_z = 26.62$
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Average pressure in lattice as function of the beam energy, estimated for three cases: a) Quadrupoles are equipped with integrated ion pumps; dipole and quadrupole chambers are glow discharge cleaned to give $\eta = 10^{-6}$ everywhere. b) and c) Quadrupoles without integrated ion pumps. These chamber sections are not glow discharge cleaned and maintain locally a high desorption yield of b) $\eta = 10^{-5}$ and c) $\eta = 10^{-4}$.

Fig. III.39
Production rate of ozone in ppm h$^{-1}$ m$^{-1}$ inside the H-magnet for a vacuum chamber with 3 mm lead screen: 
a) complete screen around vacuum chamber; 
b) lateral screens but no lead between vacuum chamber and magnet poles.
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Bunch compression in the PS by bunch shaping on the unstable fixed point: a) before phase jump of RF; b) after phase jump; c) end of bunch shaping; d) bucket re-centred on bunch; e) after rotation by 135° ejection.

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