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**THE EXPERIMENTAL DISCOVERY OF
THE INTERMEDIATE VECTOR BOSONS W^+ , W^- , AND Z^0
AT THE CERN $p\bar{p}$ COLLIDER**

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1. INTRODUCTION

1.1 The Fermi theory

It is now 50 years since Fermi [1] proposed his theory of β decay, thus solving problems which had been puzzling people for many years. In fact, with the discovery of radioactivity by Becquerel in 1896, the observation that the β rays from the decaying nucleus are electrons and that they have a continuous energy spectrum [2] led, under the assumption of a two-body final state, to the conclusion that energy and linear momentum in β -decay were not conserved.

Pauli [3] realized that also angular momentum conservation was violated, and already in 1930 he postulated the existence of a neutrino — a neutral, spin- $\frac{1}{2}$ particle with mass smaller than the electron mass. In 1934 Fermi [1], following the idea of Pauli, established the first quantitative theory of the β -decay. He assumed that the neutrino existed and that the emission of the electron-neutrino pair from the nucleus can be described in a similar way to that of photon emission from an excited atom in the radiation theory. In close analogy to quantum electrodynamics, developed several years before [4], the Lagrangian of the β -decay has the form

$$L_\beta = -(G_F/\sqrt{2})(\bar{\psi}_p\gamma^\alpha\psi_n)(\bar{\psi}_e\gamma_\alpha\psi_{\bar{\nu}_e}) \quad (1.1)$$

describing the four-fermion interaction (see Fig. 1). Here G_F is the Fermi coupling constant, determined experimentally to be

$$G_F = 1.16632 \times 10^{-5} \text{ GeV}^{-2}, \quad (\hbar = c = 1). \quad (1.2)$$

Fermi's Lagrangian is of the current-current form, its first part being the nucleon current, its second part the lepton current. It well describes many observations in β -decay, but does not account for other weak processes. As Gamow and Teller [5] pointed out, there exist β transitions which are not due to a vector interaction [Eq. (1.1)] but to an axial-vector interaction. More generally, the total interaction Lagrangian should be a linear combination of products of scalar, vector, tensor, axial-vector, and pseudoscalar terms and their Hermitian conjugates (H.c.):

$$L_\beta = -(G_F/\sqrt{2}) \sum_i C_i (\bar{\psi}_p \Omega_i \psi_n)(\bar{\psi}_e \Omega_i \psi_{\bar{\nu}_e}) + \text{H. c.}, \quad (1.3)$$

where the Ω_i are relativistic covariant combinations of γ_α , namely:

$\Omega_i = 1$	for	$i = \text{scalar (S)},$
$\Omega_i = \gamma_\alpha$	for	$i = \text{vector (V)},$
$\Omega_i = i\gamma_\alpha\gamma_\beta$	for	$i = \text{tensor (T)},$
$\Omega_i = i\gamma_\alpha\gamma_5$	for	$i = \text{axial vector (A)},$
$\Omega_i = \gamma_5$	for	$i = \text{pseudoscalar (PS)},$

and C_i are the coupling constants to be determined by experiment. In β -decay the nucleons can be treated non-relativistically, so that the first term vanishes in the case of the PS contribution.

An important step forward in the understanding of weak interactions was the appearance of the τ - θ puzzle [6], which led Lee and Yang [7] to the hypothesis that parity might be violated in weak interactions. They suggested a number of experiments that were sensitive to the question of whether

or not parity is conserved. Indeed, soon afterwards, this parity violation was observed in the decay of the polarized ^{60}Co nuclei [8] and in the decay cascade [9]:

$$\begin{aligned} \pi^\pm &\rightarrow \mu^\pm \bar{\nu}_\mu^{(-)} \\ &\quad \hookrightarrow e^\pm \bar{\nu}_e^{(-)} \bar{\nu}_\mu^{(-)}. \end{aligned} \quad (1.4)$$

Subsequent experiments on β -decay, using Mott and Møller scattering, showed that the electron is longitudinally polarized with its spin direction opposite to the direction of emission. They are “left-handed” with a helicity

$$h_{e-} = -v_e/c. \quad (1.5)$$

Positrons have

$$h_{e+} = +v_e/c. \quad (1.6)$$

To describe parity violation, Lee and Yang added to the old form of the Lagrangian [Eq. (1.3)] analogous terms which contained one and only one additional γ_5 , leading to the new form

$$L_\beta = -(G_F/\sqrt{2}) \sum_i (\bar{\psi}_p \Omega_i \psi_n) [\bar{\psi}_e \Omega_i (C_i + C'_i \gamma_5) \psi_{\bar{e}}]. \quad (1.7)$$

The terms with γ_5 are pseudoscalars and therefore parity violating. The measurements of the electron polarization showed that $C_i = C'_i$, i.e. maximum parity violation. Therefore Eq. (1.7) becomes

$$L_\beta = -(G_F/\sqrt{2}) \sum_i C_i (\bar{\psi}_p \Omega_i \psi_n) [\bar{\psi}_e \Omega_i (1 + \gamma_5) \psi_{\bar{e}}] + \text{H.c.}, \quad (1.8)$$

where the wave function of the neutrino appears with a factor $(1 + \gamma_5)$. After the discovery of parity violation it was suggested [10] that the neutrino be described by a two-component wave equation, an idea already advanced by Weyl in 1929 [11]. We consider a representation of the Dirac equation in which γ_5 is diagonal and the Dirac spinor ψ consists of two two-component spinors:

$$\psi = \begin{bmatrix} \phi \\ \chi \end{bmatrix}. \quad (1.9)$$

In this case ϕ and χ are eigenstates of γ_5 :

$$\gamma_5 \phi = \phi \quad \text{and} \quad \gamma_5 \chi = -\chi. \quad (1.10)$$

Obviously the projection operators $a^\pm = \frac{1}{2}(1 \pm \gamma_5)$ reproduce the eigenstates of γ_5 :

$$a^+ \psi = \frac{1}{2}(1 + \gamma_5) \psi = \phi \quad \text{and} \quad a^- \psi = \frac{1}{2}(1 - \gamma_5) \psi = \chi. \quad (1.11)$$

The factor $(1 + \gamma_5)$ in front of the neutrino wave function therefore means that the neutrino appears in the β -decay with a definite “chirality” $\langle \gamma_5 \rangle = +1$. Its physical meaning becomes clear using the Dirac equation, which in the case of the neutrino ($m_\nu = 0$) decouples into two independent Weyl equations:

$$[E + (\vec{p} \cdot \vec{\sigma})]\phi = 0 \quad \text{and} \quad [E - (\vec{p} \cdot \vec{\sigma})]\chi = 0. \quad (1.12)$$

Here E and $\vec{p}$ are the energy and momentum, with $|\vec{p}| = E$. From Eqs. (1.12) it follows immediately that

$$(\vec{p} \cdot \vec{\sigma})/|\vec{p}| = h = -1 \text{ for } \phi \quad \text{and} \quad h = +1 \text{ for } \chi. \quad (1.13)$$

Therefore, representing the neutrino and antineutrino by ϕ and χ , respectively, Eqs. (1.13) show that neutrinos are always left-handed and antineutrinos always right-handed:

$$\langle \gamma_5 \rangle = -h \text{ for } \nu \quad \text{and} \quad \langle \gamma_5 \rangle = +h \text{ for } \bar{\nu}. \quad (1.14)$$

We now have to find an interaction which produces only left-handed neutrinos, since ψ_ν appears only in combination with $(1 + \gamma_5)$. This factor has consequences for the electron wave function. It can be shown that in the (V,A) case, there must be a factor $\frac{1}{2}(1 + \gamma_5)$ in front of $\bar{\psi}_e$, whereas a factor $\frac{1}{2}(1 - \gamma_5)$ should be introduced in the (S,T,P) case. In other words, electrons and neutrinos have the same chirality for (V,A), and opposite for (S,T,P). Therefore we need to measure the relative helicity of electrons and antineutrinos.

But already before these measurements it was postulated [12] that requiring invariance under the transformation

$$\psi \rightarrow \gamma_5 \psi \quad (1.15)$$

for all wave functions would produce the right interaction. After the discovery of parity violation, this γ_5 invariance was reformulated [13] to

$$\psi \rightarrow [(1 + \gamma_5)/2]\psi = \phi \quad (1.16)$$

for all ψ , i.e. each fermion enters with definite chirality and as a two-component spinor. To obtain the true Lagrangian for the four-fermion interaction, we just need to replace in Eq. (1.3) all ψ by $\phi = [(1 + \gamma_5)/2]\psi = a^+ \psi$:

$$L_\beta = -\sqrt{8G_F} \sum_i C_i (\bar{a}^+ \psi_p \Omega_i a^+ \psi_n) (\bar{a}^+ \psi_e \Omega_i a^+ \psi_{\bar{\nu}_e}) + \text{H.c.}, \quad (1.17)$$

or

$$L_\beta = -\sqrt{8G_F} \sum_i C_i (\bar{\psi}_p a^- \Omega_i a^+ \psi_n) (\bar{\psi}_e a^- \Omega_i a^+ \psi_{\bar{\nu}_e}) + \text{H.c.} \quad (1.18)$$

It now turns out that

$$a^- \Omega_i = 0 \text{ for S,T,P} \quad \text{and} \quad a^- \Omega_i \neq 0 \text{ for V,A.} \quad (1.19)$$

Expressing Ω_V and Ω_A by their γ matrices, Eq. (1.18) becomes then

$$L_\beta = -(G_F/\sqrt{2}) [\bar{\psi}_p \gamma^\alpha (C_V - C_A \gamma_5) \psi_n] [\bar{\psi}_e \gamma_\alpha (1 + \gamma_5) \psi_{\bar{\nu}_e}] + \text{H.c.} \quad (1.20)$$

But the experimental results have to decide upon this V,A theory. The first step was the measurement of the relative helicity between electrons and antineutrinos in nuclear β -decay. After many controversial results, the experiments on ^{35}A , ^{19}Ne , ^{23}Ne , and ^6He [14] indicated that V,A coupling is dominant. The same result was obtained by measuring the helicity of the neutrino directly through the circular polarization of γ -rays from ^{152}Sm [15]. The result is

$$h_{\nu_e} = -1. \quad (1.21)$$

Only the phase between V and A has to be determined. This implies an interference experiment between V and A couplings, an effect which is automatically parity violating. The experiment [16] measured the angular distribution of the electrons from polarized neutrons. From the resulting asymmetry coefficients for electrons and neutrinos it turned out that the phase is 180° , leading to $C_V \approx -C_A$. Therefore we are speaking of a V - A interaction. The exact value for C_V was determined in a pure vector decay ($^{14}\text{O} \rightarrow ^{14}\text{N}$), the value for C_A from the measured lifetime of the neutron. The present ratio [17]

$$\lambda = -C_A/C_V = 1.25 \pm 0.009 \quad (1.22)$$

deviates slightly from V - A, but this is attributed to the strong interaction of the nucleons. The final Lagrangian then becomes

$$L_\beta = -(G_F/\sqrt{2}) [\bar{\psi}_p \gamma^\alpha (1 + \lambda \gamma_5) \psi_n] [\bar{\psi}_e \gamma_\alpha (1 + \gamma_5) \psi_{\bar{\nu}_e}] + \text{H.c.} \quad (1.23)$$

and, apart from the important γ_5 factor, it is Fermi's first "ansatz" [Eq. (1.1)], but it took 25 years to discover $(1 + \gamma_5)$.

Equation (1.23) suggests that the four-fermion interaction is described by a single universal current J^α . This results in the "universal Fermi interaction" [13], where the effective Lagrangian is described by a coupling between the total weak current J^α and its Hermitian conjugate:

$$L = -(G_F/\sqrt{2}) J^\alpha J_\alpha^\dagger. \quad (1.24)$$

The components of J^α are the well-known particle currents:

the lepton currents

$$\ell_e^\alpha = [\bar{\psi}_e \gamma^\alpha (1 + \gamma_5) \psi_{\nu_e}] \quad \text{and} \quad \ell_\mu^\alpha = [\bar{\psi}_\mu \gamma^\alpha (1 + \gamma_5) \psi_{\nu_\mu}] \quad (1.25)$$

and the hadron current

$$h^\alpha = [\bar{\psi}_n \gamma^\alpha (1 + \lambda \gamma_5) \psi_p]. \quad (1.26)$$

Equation (1.24) contains all reactions which we already know. In addition, there now appear reactions that were experimentally unknown. In fact the known reactions are products between ℓ_1^α and h^α , the so-called “non-diagonal” terms. If the idea of the universal current–current coupling is exact, then there appear “diagonal” terms — where each current couples with itself — such as $\ell_e^\alpha \ell_e^\dagger$ or $h^\alpha h^\dagger$ etc., resulting in $\nu_e e^-$ or np scattering etc. These reactions were indeed observed.

1.2 The difficulty of the Fermi interaction

and the need for intermediate vector bosons

Although the universal Fermi interaction [Eq. (1.24)] provides an excellent description of the weak interaction processes, it cannot be concealed that fundamental difficulties exist, a fact already recognized soon after Fermi’s postulation [18]. Consider the process

$$\nu_e e^- \rightarrow \nu_e e^- \quad (1.27)$$

mentioned above. The total cross-section is

$$\sigma(\nu_e e^-) = 2G_F^2 m_e E_{\nu_e} / \pi = G_F^2 s / \pi, \quad (1.28)$$

where $s = 2m_e E_{\nu_e}$ is the square of the centre-of-mass energy E_{cm} . In the wave theory, the total cross-section may be written as a sum of partial wave cross-sections,

$$\sigma = (\pi p_{\text{cm}}^2) \sum_J (2J+1) |f_J|^2, \quad (1.29)$$

where f_J is the amplitude of the J^{th} partial wave, and p_{cm} is the centre-of-mass momentum of the particles ($p_{\text{cm}} = \sqrt{s}/2$). A point-like interaction requires $J = 0$, and as a consequence of unitarity (= conservation of probability) it yields $|f_J| \leq 1$. Therefore Eq. (1.29) becomes

$$\sigma \leq 4\pi/s. \quad (1.30)$$

Comparing it with Eq. (1.28) it must be that

$$G_F^2 s / \pi \leq 4/\pi s. \quad (1.31)$$

The unitarity bound is therefore violated for

$$p_{\text{cm}} = \sqrt{s}/2 \leq (\pi/2G_F)^{1/2} \approx 367 \text{ GeV}. \quad (1.32)$$

At sufficiently large energy, the Fermi theory of weak interaction is therefore inconsistent. Also, inclusion of higher-order corrections does not remove the divergences inherent to the theory. The basic problem lies in the fact that G_F is not dimensionless; it has dimension $[E^{-2}]$ [see Eq. (1.28)]. If we look at an analogous process in quantum electrodynamics (QED), e.g. $e^+e^- \rightarrow \mu^+\mu^-$, we find that the total cross-section for this process is

$$\sigma \sim \alpha^2/E_{\text{cm}}^2 = e^4/E^2, \quad (1.33)$$

where α , the coupling constant, is dimensionless, so that E_{cm}^2 is in the denominator to give σ the dimension $[E^{-2}]$. The cross-section (1.33) falls with increasing energy and will not violate unitarity. QED is a renormalizable theory. Divergences in higher-order diagrams are eliminated by mass and charge renormalization, a method which fails in the weak interaction theory.

It is therefore very appealing to redefine the $V - A$ theory in terms of a dimensionless coupling constant. Hence we reject the idea of a point-like interaction and assume that the weak interaction proceeds via the exchange of an intermediate particle [19], like the photon in QED and in analogy to Yukawa's meson theory of nuclear forces [20]. This particle must have the following characteristics: the spin = 1 because of $V - A$; the charge = ± 1 because we observe charged currents; and the mass $\neq 0$ because of the short range of the weak interaction. We call this new particle the intermediate vector boson (IVB); in fact there are two: W^+ and W^- . The well-known β -decay now proceeds differently (Fig. 2) with a new coupling constant g_w at each vertex. This new non-local interaction automatically introduces a new term in the Lagrangian (1.24) of the form

$$g_w^2/(q^2 + m_w^2), \quad (1.34)$$

where the denominator is the propagator, with q^2 being the four-momentum transfer. For $m_w^2 \gg q^2$ or $q^2 \rightarrow 0$ we have to find again the old universal Fermi interaction (1.24), successfully describing the experiments at low energy. Therefore

$$G_F/\sqrt{2} = \lim_{q^2 \rightarrow 0} g_w^2/(q^2 + m_w^2) \quad (1.35)$$

or

$$G_F/\sqrt{2} = g_w^2/m_w^2. \quad (1.36)$$

The first neutrino experiments at accelerators could only arrive at a lower limit of the boson mass, ranging from 1.7 to 2.2 GeV/c² depending on the leptonic branching ratio [21].

The Lagrangian (1.24) now becomes

$$L = -g_w(J^\alpha W_\alpha^\dagger + J_\alpha^\dagger W^\alpha), \quad (1.37)$$

where W_α is the field of the IVB, and J_α is the total weak current, which is the sum of lepton and hadron currents,

$$J^\alpha = \ell^\alpha + h^\alpha. \quad (1.38)$$

This important modification to the universal Fermi interaction is still not satisfactory. In fact the unitarity limit is still violated for large energies in a given partial-wave amplitude, although logarithmically. Furthermore, higher-order terms diverge again quadratically in E_{cm} , leading to $\sigma \sim s$, e.g. in the process $\nu_e \bar{\nu}_e \rightarrow W^+ W^-$ (see Fig. 3a). Therefore the theory is still not renormalizable.

The idea is now to compensate for the divergences in the diagram of Fig. 3a by adding other diagrams so that the sum of all diagrams yields a cross-section which does not violate unitarity. One possibility is shown in Fig. 3b. It leads to the introduction of a neutral intermediate boson Z^0 which has the same characteristics as the $W^\pm$, except for the charge. Its mass should be of order m_W . This seemed to be a plausible solution [22, 23] to the above difficulties. It now turns out that it is possible to construct a theory which takes the latter assumption into account and makes sensible predictions at high energies and also for higher-order processes: the gauge theory of electroweak interaction. It is renormalizable, a fact which in QED is strongly coupled to gauge invariance.

2. THE STANDARD MODEL

A simple model for a gauge theory of electroweak interaction, the so-called standard model, has been written down originally by Glashow [23], and in its final version by Weinberg and Salam [24]. Here we may only describe the important lines of this theory; for more details we refer the reader to the literature [25]. The standard model is based essentially on the ideas of generalized gauge invariance, spontaneous symmetry breaking, and the Higgs mechanism.

Let us start with the interaction of electrons and photons. The Lagrangian L_{em} of QED states an important fact: the electromagnetic current is locally conserved. This is due to the invariance of L_{em} under a local gauge transformation. For this local gauge invariance the existence of the photon is essential. Therefore the photon field A_μ is an example of a gauge field. In addition, the gauge invariance of L_{em} ensures the masslessness of the photon.

We now use the form of L_{em} to construct a Lagrangian L_W which describes the weak interaction of the intermediate boson field W_μ with the leptons. The problem with such a L_W is that it is not gauge invariant at all and that such a theory is not renormalizable. A possible solution is given by the Yang-Mills model [26]. There, an isotriplet of massless spin-1 particles (W^+ , W^0 , W^-) interacts, for instance, with a fermion isodoublet in such a way that isospin is conserved. The corresponding Lagrangian L_{YM} has the characteristics that it leads to the local conservation of the isospin current. This is again due, as in QED, to the invariance of L_{YM} under a special gauge transformation, showing that the isovector field $\vec{W}$ is a gauge field. This looks very promising for constructing a theory of weak interaction. We give to the fermions the hypercharge Y which couples gauge-invariantly to a massless spin-1 field B_μ . Together with the isospin I the fermions then couple in a *generalized gauge-invariant* way to the fields $\vec{W}_\mu$ and B_μ . The total Lagrangian will now possess $SU(2) \times U(1)$ invariance. But all the particles — fermions and bosons — still have zero mass. We have to introduce another particle, a spin-0 boson carrying isospin and hypercharge, in order that the particles acquire a mass.

Before describing how this happens, we first consider the spin-0 field in isolation to illustrate an effect known as *spontaneous symmetry breaking*. It can be shown that for a scalar quantum field ϕ with $I = 0$, $Y = 0$, having some self-interaction $V(\phi)$, the ground state is the vacuum expectation value

$$\langle 0 | \phi | 0 \rangle \equiv \eta \neq 0 \quad (2.1)$$

for a special choice of $V(\phi)$. A Lagrangian in terms of $\chi = \phi - \eta$ can be interpreted as describing the spin-0 field χ with a mass $\sim 2\eta$. Thus, choosing $V(\phi)$ so that $\langle 0 | \phi | 0 \rangle \neq 0$, the spin-0 particle appears to acquire a mass.

Next, giving to the particle a hypercharge, we have two real fields ϕ_1 and ϕ_2 . The same result is now obtained, with $V(\phi, \phi^*)$, but in addition it follows that the vacuum state $|0\rangle$ is not an eigenstate of hypercharge, whereas the Lagrangian for such a system is invariant under the hypercharge gauge transformation. This is known as spontaneous breaking of hypercharge gauge invariance. The mass for χ_1 is again $\sim 2\eta$, whereas the mass for $\chi_2 = 0$. So the Lagrangian describing two fields ϕ_1 and ϕ_2 in a symmetric way, now describes two scalar particles of different mass: symmetry is broken. The massless particle is called the Goldstone boson [27].

If finally we look at an isovector real spin-0 field ($I = 1, Y = 0$), $\vec{\phi} = (\phi_1, \phi_2, \phi_3)$, we will find that the Lagrangian conserves isospin but that the vacuum is not an eigenstate of isospin. We have spontaneously broken isospin symmetry. In the mass spectrum we find two massless particles (Goldstone bosons) and one with mass. If we let such a spin-0 particle interact with a gauge field, we get the *Higgs effect* [28].

We consider, as above, a complex field ϕ with hypercharge which can be expressed in terms of two real fields θ and σ . Because of the special form of $V(\phi, \phi^*)$, the vacuum expectation value has a non-zero value and is chosen to be

$$\langle 0 | \phi | 0 \rangle = \eta = \text{real}. \quad (2.2)$$

The field ϕ now interacts with a massless gauge field B_μ , leading to a gauge-invariant coupling with a coupling constant g' . The complete Lagrangian, after some substitutions, has the following properties: first it describes a spin-1 boson B_μ with a non-zero mass $g'\eta$; this is a result of the interaction with ϕ . Next, it appears that the scalar field σ has also a mass [due to $V(\sigma)$] but the θ -field has completely disappeared, i.e. the Goldstone boson, corresponding to the θ -field, has been absorbed by the gauge field B_μ , which is now massive. The non-appearance of the Goldstone boson and the occurrence of the massive spin-1 particle is known as the Higgs-Kibble mechanism [28, 29].

We are now in a position to construct the standard model. As described before, we start with four types of particles, all of which have zero mass:

- i) leptons ($s = \frac{1}{2}$) ℓ ,
- ii) quarks ($s = \frac{1}{2}$) q ,
- iii) vector bosons ($s = 1$) W, B ,
- iv) a scalar boson ($s = 0$) ϕ .

We assume that the leptons and quarks occur in $SU(2)$ multiplets, i.e. they carry a weak isospin I . In addition they possess hypercharge Y connected to the electric charge via

$$Q = I_3 + Y, \quad (2.3)$$

where I_3 is the third component of the weak isospin. The vector bosons consist of an isotriplet $\vec{W}$ and an isosinglet B , i.e. they have also weak isospin I , but they carry no hypercharge Y . The scalar boson ϕ has $I = \frac{1}{2}$ and $Y = \frac{1}{2}$, which through [Eq. (2.3)] leads to

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}. \quad (2.4)$$

Next we consider the different interactions between the particles:

- 1) The isovector field $\vec{W}_\mu$ couples gauge-invariantly with g to all particles carrying weak isospin I , i.e. to ℓ , q , and ϕ . The Lagrangian is $SU(2)$ gauge-invariant.
- 2) The isoscalar field B_μ couples gauge-invariantly with g' to all particles carrying hypercharge, i.e. to ℓ , q , and ϕ . The Lagrangian is therefore $U(1)$ gauge-invariant.
- 3) The scalar field ϕ has Yukawa couplings with ℓ and q , conserving I and Y .
- 4) The scalar field has in addition a self-interaction $V(\phi)$.

From interactions (1) and (2) it follows that the corresponding Lagrangian is $SU(2) \times U(1)$ invariant. It is this group that characterizes the standard model. So far all particles have been massless. If we assume that $V(\phi)$ has a form such that $\langle 0 | \phi | 0 \rangle = \eta \neq 0$, the vacuum is not an eigenstate of I and Y . Because of the Higgs effect, three of the four vector bosons acquire mass, the fourth one remaining massless, leading to the following interpretation:

fields $W_\mu^\pm$ ($m_{W^\pm}^2 = 1/4 g^2 \eta^2$) $\hat{=}$ charged intermediate bosons W^+ and W^- ;

field W_μ^0 ($m_{W^0}^2 = 1/4 \eta^2 (g^2 + g'^2)$) $\hat{=}$ neutral intermediate boson Z^0 ; (2.5)

field A_μ ($m_A^2 = 0$) $\hat{=}$ photon γ .

Also, the leptons and quarks now obtain a mass because of their couplings to ϕ , but the neutrino remains massless. From the original ϕ there is left only one single massive neutral scalar particle H^0 , the Higgs boson.

Let us turn to the interesting interactions between W_μ, B_μ and ℓ, q . The interaction Lagrangian is

$$L_{ew} = g \vec{J}_\mu \vec{W}_\mu + g' J_\mu^{(Y)} B_\mu, \quad (2.6)$$

where $\vec{J}_\mu$ is the isospin current and $J_\mu^{(Y)}$ the hypercharge current; g and g' are the coupling constants. Because $Y = Q - I_3$, we can write

$$J_\mu^{(Y)} = J_\mu^{(e)} - J_\mu^{(3)}, \quad (2.7)$$

where $J_\mu^{(e)}$ is the electromagnetic current, and we obtain

$$L_{ew} = g \vec{J}_\mu \vec{W}_\mu + g' [J_\mu^{(e)} - J_\mu^{(3)}] B_\mu. \quad (2.8)$$

Expressing the fields $\vec{W}_\mu, B_\mu$ in terms of the $W_\mu^+, W_\mu^-, Z_\mu^0$, and A_μ we find

$$\begin{aligned} W_\mu^\pm &= [W_\mu^{(1)} \pm i W_\mu^{(2)}] / \sqrt{2} \\ Z_\mu^0 &= [g W_\mu^{(3)} + g' B_\mu] / (g^2 + g'^2)^{1/2} \\ A_\mu &= [-g' W_\mu^{(3)} + g B_\mu] / (g^2 + g'^2)^{1/2}. \end{aligned} \quad (2.9)$$

Introducing the weak angle θ_w so that

$$\tan \theta_w = g'/g, \quad (2.10)$$

the Lagrangian (2.8) becomes

$$L_{ew} = (g/2\sqrt{2}) [2J_\mu^- W_\mu^+ + 2J_\mu^+ W_\mu^-] + (g/\cos \theta_w) [J_\mu^{(3)} - J_\mu^{(e)} \sin^2 \theta_w] Z_\mu + g' J_\mu^{(e)} A_\mu \cos \theta_w, \quad (2.11)$$

with $J_\mu^{(\pm)} \equiv J_\mu^{(1)} \pm iJ_\mu^{(2)}$.

The first term is the well-known weak interaction term (1.37); the third term is the standard electromagnetic interaction. The second term is new. It describes neutral currents in weak interactions, a phenomenon not observed at the time when the standard model was postulated.

It remains for us to classify the leptons and quarks. The choices are very much constrained by the requirement that the new theory must reproduce the results of QED and the $V - A$ character of the weak interaction. Therefore the leptons are grouped into left-handed (L) doublets and right-handed (R) singlets:

$$\begin{bmatrix} \nu_e \\ e \end{bmatrix}_L, \quad \begin{bmatrix} \nu_\mu \\ \mu \end{bmatrix}_L, \quad \begin{bmatrix} \nu_\tau \\ \tau \end{bmatrix}_L, \quad I = 1/2, Y = -1/2, \quad (2.12)$$

and

$$e_R, \mu_R, \tau_R, \quad I = 0, Y = 1. \quad (2.13)$$

The same is true for the quarks:

$$\begin{bmatrix} u \\ d_c \end{bmatrix}_L, \quad \begin{bmatrix} c \\ s_c \end{bmatrix}_L, \quad \begin{bmatrix} t \\ b_c \end{bmatrix}_L, \quad I = 1/2, Y = 1/6, \quad (2.14)$$

and

$$u_R, d_{cR}, c_R, s_{cR}, t_R, b_{cR}, \quad I = 0, Y = Q, \quad (2.15)$$

where the index c indicates that we deal with Cabibbo–GIM–KM-rotated quarks [30]. The $V - A$ law is thus inserted “by hand” in the standard model. The appearance of three lepton and three quark families is due to the discovery, in the last ten years, of the τ lepton [31], the charmed quark c [32], and the bottom (or beauty) quark b [33]. The top (or truth) quark t still remains undiscovered, but there seem to be indications of its existence. In fact, the axial-vector coupling constant of the b quark has been determined to be $g_A^b = -0.45 \pm 0.25$ [34], in agreement with -0.5 expected for $(\frac{t}{b_c})_L$. As a result, we have to add the leptonic current ℓ_τ^α in Eq. (1.25), and can decompose (1.26) into the individual quark currents.

We were left with the question of renormalization. This difficult problem was solved in 1971 and it turned out that a massless Yang–Mills theory is indeed renormalizable [35] and that the same is true for the standard model [36].

Finally, let us exploit the Lagrangian (2.11). The electromagnetic interaction has the coupling constant e . Therefore, comparing it with the third term in Eq. (2.11), it must be

$$g \sin \theta_w = g' \cos \theta_w = e. \quad (2.16)$$

This is the unification condition. It means that the coupling to $W^\pm$ and Z^0 is of the same order of magnitude as the electromagnetic coupling e .

From the first term in Eq. (2.11), and comparing it with Eq. (1.37), we end up with [see Eq. (1.36)]

$$g/2\sqrt{2} = g_W = m_W(G_F/\sqrt{2})^{1/2} \quad (2.17)$$

or

$$m_W = (e^2\sqrt{2}/8G_F)^{1/2} / \sin \theta_w (\text{GeV}/c^2). \quad (2.18)$$

Using the mass relation in Eq. (2.5) we find

$$m_W^2 = m_Z^2 \cos^2 \theta_w. \quad (2.19)$$

A new parameter, ρ , was introduced [37] in the theory, which is directly related to the Higgs system by

$$\rho = m_W^2/m_Z^2 \cos^2 \theta_w = 1/2 I_\phi, \quad (2.20)$$

where I_ϕ is the isospin of the scalar boson ϕ . In the standard model, $\rho = 1$.

In conclusion we can say that the standard model reproduces the phenomenology of electromagnetic and weak interactions. It gives definite values for masses of the intermediate bosons $W^\pm$ and Z^0 and has only one parameter, the mixing angle θ_w . Furthermore, it predicts the existence of neutral currents in weak interactions, e.g. the pure leptonic elastic scattering process (Fig. 4),

$$\nu_\mu(\bar{\nu}_\mu)e^- \rightarrow \nu_\mu(\bar{\nu}_\mu)e^-, \quad (2.21)$$

a process previously inconceivable. The great experimental activity finally discovered [38] process (2.21), and soon afterwards the neutral currents in neutrino-quark reactions [39]. To date, neutral-current phenomena have been observed in ν reactions, in e and μ reactions, in e^+e^- reactions, in atomic physics, etc. [34, 40, 41]. They all confirm the standard model with a value [41] for

$$\sin^2 \theta_w = 0.233 \pm 0.009. \quad (2.22)$$

Although the standard model is very successful, it was not until 1983 that the intermediate bosons $W^\pm$ and Z^0 were discovered [42–46], 25 years after $(1 + \gamma_s)$.

3. THE GENERAL PROPERTIES OF THE INTERMEDIATE VECTOR BOSONS

As we have seen, the standard model makes precise predictions about the existence of three intermediate bosons, W^+ , W^- , and Z^0 , and especially about their mass values, mainly depending on

the only free parameter of the theory, the weak mixing angle θ_w . In the following, we discuss the properties [47–50] of the IVBs in relation to the experiments that led to their discovery.

3.1 Properties of the W^+ and W^-

With the substitution $e^2 = 4\pi\alpha$, $\alpha = 1/137.036$ being the fine-structure constant, we find for Eq. (2.18)

$$m_W = (\pi\alpha/\sqrt{2} G_F)^{1/2} / \sin \theta_w = 37.3/\sin \theta_w \text{ GeV}/c^2. \quad (3.1)$$

Using the value of $\sin^2 \theta_w$ from Eq. (2.22), this gives

$$m_W = 77.2 \pm 1.6 \text{ GeV}/c^2. \quad (3.2)$$

However, the above relations are of lowest order. In fact, sizeable $O(\alpha)$ radiative corrections affect the values. As an example, the fine structure constant α becomes $\alpha(m_W) = 1/127.7$ at mass scale m_W . Since the theory is renormalizable, these corrections have been calculated [51, 52]. Following Ref. 52, we find

$$m_W = 38.5/\sin \theta_w^R \text{ GeV}/c^2, \quad (3.3)$$

where $\sin \theta_w^R$ is now the renormalized value at mass scale m_W , defined by minimal subtraction, $\overline{MS}$, in the dimensional regularization scheme, with a value of

$$\sin^2 \theta_w^R = 0.215 \pm 0.014, \quad (3.4)$$

leading to the present prediction [52]

$$m_W = 83.0_{-2.8}^{+3.0} \text{ GeV}/c^2. \quad (3.5)$$

Creating particles with such a large mass remained a daunting task until the idea of a $p\bar{p}$ collider was considered.

Next we are interested in the different decay modes of the W 's, e.g. for the W^- :

$$\begin{aligned} W^- &\rightarrow \ell \bar{\nu}_\ell & \ell \bar{\nu}_\ell &= (e^- \bar{\nu}_e), (\mu^- \bar{\nu}_\mu), (\tau^- \bar{\nu}_\tau) \\ W^- &\rightarrow q \bar{q}' & q \bar{q}' &= (d\bar{u}), (s\bar{c}), (b\bar{t}). \end{aligned} \quad (3.6)$$

For the W^+ we have to replace the particle by its antiparticle and vice versa. From the first term in (2.11), which describes the coupling between the weak current (leptonic and hadronic) and the charged intermediate bosons, we can calculate the partial decay width for all six modes in Eq. (3.6). We find

$$\Gamma(W \rightarrow \ell \nu_\ell) = G_F m_W^3 / 6\pi\sqrt{2} = 250 \text{ MeV} \quad (3.7)$$

using the value of m_W from (3.5).

The partial rates for the decay into quark–antiquark pairs are, in first order, three times larger than in Eq. (3.7) owing to the colour the quarks carry and because all m_q/m_W effects are neglected. This is admissible except for the unseen top quark, whose mass is limited to $m_t \geq 22.5 \text{ GeV}/c^2$ from PETRA results [53]. Taking this into account and including QCD corrections [52], we find for the sum of all three $q\bar{q}'$ modes in (3.6),

$$\begin{aligned}\Gamma(W \rightarrow \text{hadrons}) &= (G_F m_W^3 / 2\pi\sqrt{2}) [1 - (m_t^2/2m_W^2) + (m_t^6/6m_W^6)] (1 + \alpha_s^R/\pi) \\ &= 2.20 \text{ GeV} \quad (\text{for } m_t = 30 \text{ GeV}/c^2),\end{aligned}\tag{3.8}$$

where α_s^R is the strong coupling constant at mass scale m_W . With Eq. (3.7), the total decay width is therefore

$$\Gamma(W \rightarrow \text{all}) = 2.95 \text{ GeV}.\tag{3.9}$$

This leads to a leptonic branching ratio—interesting in the $e^\pm$ and the $\mu^\pm$ cases from the experimental point of view—of

$$\Gamma(W \rightarrow \ell\nu)/\Gamma(W \rightarrow \text{all}) = 8.5\%.\tag{3.10}$$

3.2 Properties of the Z^0

Using the mass formulas (2.19) and (3.3) as well as the $\sin^2 \theta_W^R$ value (3.4), the prediction for the Z^0 mass [52] is

$$m_{Z^0} = 93.8_{-2.4}^{+2.5} \text{ GeV}/c^2.\tag{3.11}$$

Since the Z^0 is neutral, the number of decay modes is now much larger, namely

$$\begin{aligned}Z^0 \rightarrow \ell\bar{\ell} \quad \ell\bar{\ell} &= (e^-e^+), (\nu_e\bar{\nu}_e), (\mu^-\mu^+), (\nu_\mu\bar{\nu}_\mu), (\tau^-\tau^+), (\nu_\tau\bar{\nu}_\tau); \\ Z^0 \rightarrow q\bar{q} \quad q\bar{q} &= (d\bar{d}), (u\bar{u}), (s\bar{s}), (c\bar{c}), (b\bar{b}), (t\bar{t}).\end{aligned}\tag{3.12}$$

Now the second term in Eq. (2.11) delivers the partial decay width for all the above decay modes. In general we have

$$\Gamma(Z^0 \rightarrow \ell\bar{\ell}) = (G_F m_Z^3 / 6\pi\sqrt{2}) (g_V^2 + g_A^2)\tag{3.13}$$

and

$$\Gamma(Z^0 \rightarrow q\bar{q}) = (G_F m_Z^3 / 2\pi\sqrt{2}) (q_V^2 + q_A^2).\tag{3.14}$$

Note that the constant term in Eq. (3.14) differs again from Eq. (3.13) by the colour factor 3. The parameters g_V, q_V and g_A, q_A are the chiral, vector, and axial-vector coupling constants^{*)} between the Z^0 and the leptons, respectively the quarks:

^{*)} Sometimes in the literature one finds c_V and c_A : $c_V = g_V$ and $c_A = g_A$ and the q_V and q_A are normally expressed in terms of the different quarks, like u_V, u_A, d_V, d_A , etc.

$$\begin{aligned} g_V &= I_3^\ell - 2|Q_\ell| \sin^2 \theta_w; & q_V &= I_3^q - 2|Q_q| \sin^2 \theta_w \\ g_A &= I_3^\ell & q_A &= I_3^q, \end{aligned} \quad (3.15)$$

where I_3^ℓ , I_3^q are the third components of the weak isospin, and Q_ℓ , Q_q are the electric charges for both leptons (ℓ) and quarks (q). For the leptons we find the following partial decay widths:

$$\begin{aligned} \Gamma(Z^0 \rightarrow \ell^+ \ell^-) &= 92 \text{ MeV} & \text{with } \ell &= e, \mu, \tau, \\ \Gamma(Z^0 \rightarrow \nu \bar{\nu}) &= 181 \text{ MeV} & \text{with } \nu &= \nu_e, \nu_\mu, \nu_\tau. \end{aligned} \quad (3.16)$$

As for the $W^\pm$, we can again calculate the total decay width into hadrons, taking the heaviness of the top quark into account. We find, for $m_t = 30 \text{ GeV}/c^2$,

$$\Gamma(Z^0 \rightarrow \text{hadrons}) = 2.12 \text{ GeV}, \quad (3.17)$$

so that the total width becomes

$$\Gamma(Z^0 \rightarrow \text{all}) = 2.94 \text{ GeV} \quad (3.18)$$

leading to a branching ratio of Z^0 into charged dileptons of

$$\Gamma(Z^0 \rightarrow \ell^+ \ell^-) / \Gamma(Z^0 \rightarrow \text{all}) = 3.1\%. \quad (3.19)$$

It is interesting to note that the total decay width is sensitive to the number of neutrinos. Since we do not know the number of families involved in particle interactions, and since we believe that neutrinos are much lighter than their corresponding charged lepton partner, the Z^0 can decay always into neutrino pairs, whereas the $W^\pm$ cannot decay into a lepton and its neutrino if the mass of the lepton approaches m_W . So a precise measurement of the Z^0 width gives a direct answer to the number of families, since each new family adds 181 MeV to the total width through the $Z^0 \rightarrow \nu \bar{\nu}$ decay.

3.3 The production of $W^\pm$ and Z^0

Massive particles, such as $W^\pm$ and Z^0 , can be produced in lepton–lepton, lepton–hadron, or hadron–hadron interactions, provided the available energy is large enough to produce them. We will concentrate on the hadron–hadron interactions, say $p\bar{p}$ collisions, which led finally to the discovery of the IVBs.

For the production of the IVBs in $p\bar{p}$ collisions one uses as a guide the Drell–Yan process [54] in the framework of the quark–parton model. This process was originally used to describe lepton–pair production in hadron–hadron collisions. In this model, a quark from one hadron annihilates with an antiquark from the other, producing a virtual photon which in turn produces a lepton pair (Fig. 5a). In this picture quarks and antiquarks are supposed to be free objects inside the hadrons [55], so that the differential cross-section in terms of the mass M and the momentum fraction x of the lepton pair becomes [48]

$$d^2\sigma/dM^2 dx = (4\pi\alpha^2/3M^4) F(\tau, x), \quad (3.20)$$

where the factor in parenthesis reflects the point-like character of the scattering. The function $F(\tau, x)$ describes all processes involved in the interaction. It depends on the dimensionless scaling variable

$$\tau = M^2/s = x_1 x_2, \quad (3.21)$$

where x_1, x_2 are the momentum fractions carried by the constituents related to x :

$$|x_1 - x_2| = x = 2p_L/\sqrt{s}, \quad (3.22)$$

p_L being the longitudinal momentum of the lepton pair. Formula (3.21) is only valid if one neglects the momenta of the constituents transverse to the beam direction. The function $F(\tau, x)$ is the sum over all reactions:

$$F(\tau, x) = [x_1 x_2 / (x^2 + 4\tau)^{1/2}] R(x_1, x_2), \quad (3.23)$$

with

$$R(x_1, x_2) = (1/3) \sum_i e_i^2 [q_i(x_1) \bar{q}_i(x_2) + \bar{q}_i(x_1) q_i(x_2)], \quad (3.24)$$

where e_i is the charge of the quark flavour i , $q_i(x_1)$ is the probability of finding quark i with momentum fraction x_1 inside the hadron 1, $\bar{q}_i(x_2)$ the probability of finding antiquark i with x_2 in hadron 2, and so on. The factor $1/3$ in front corresponds to colour averaging. The integrated cross-section becomes after substitution

$$d\sigma/dM^2 = (4\pi\alpha^2/3M^4) \tau \int_{\tau}^1 dx R(x, \tau/x)/x \quad (3.25)$$

This implies that $M^4 d\sigma/dM^2$ should scale since it depends only on τ . Indeed, in recent years, the Drell–Yan model has been tested experimentally; it gives, in fact, a good description of the data, and scaling is found to be reasonably well satisfied up to $\sqrt{\tau} = 0.5$ [56].

The success of this naïve model therefore suggests its application also for the coupling of the weak currents to the $W^\pm$ - and Z^0 -production [47–50, 57, 58] (Fig. 5b). For the $W^\pm$ -production one finds, with the relevant elementary processes, for instance:

$$u + \bar{d} \rightarrow W^+, \quad u + \bar{s} \rightarrow W^+, \quad (3.26)$$

$$\bar{u} + d \rightarrow W^-, \quad \bar{u} + s \rightarrow W^-,$$

a differential cross-section

$$d\sigma/dx = \sqrt{2} G_F \pi [x_1 x_2 / (x^2 + 4\tau)^{1/2}] w^\pm(x_1, x_2). \quad (3.27)$$

Here the quark–quark interaction in the case of the W^+ is described by

$$\begin{aligned} w^+(x_1, x_2) = 1/3 \{ & [u_1(x_1) \bar{d}_2(x_2) + \bar{d}_1(x_1) u_2(x_2)] \cos^2 \theta_c \\ & + [u_1(x_1) \bar{s}_2(x_2) + \bar{s}_1(x_1) u_2(x_2)] \sin^2 \theta_c \} \end{aligned} \quad (3.28)$$

where θ_c is the Cabibbo angle.

For W^- -production one has to interchange the quark and antiquark densities. The total cross-section becomes then

$$\sigma_{W^\pm} = \sqrt{2} G_F \pi \tau \int_{\tau}^1 dx w^\pm(x, \tau/x)/x. \quad (3.29)$$

In the case of the Z^0 -production we have a function $z^0(x_1, x_2)$ describing for instance the processes

$$u + \bar{u} \rightarrow Z^0, \quad d + \bar{d} \rightarrow Z^0, \quad s + \bar{s} \rightarrow Z^0, \quad (3.30)$$

defined as

$$\begin{aligned} z^0(x_1, x_2) = & 1/3 \{ [u_1(x_1)\bar{u}_2(x_2) + \bar{u}_1(x_1)u_2(x_2)] [1/4 - (2/3)\sin^2 \theta_w + (8/9)\sin^4 \theta_w] \\ & + [d_1(x_1)\bar{d}_2(x_2) + \bar{d}_1(x_1)d_2(x_2) + s_1(x_1)\bar{s}_2(x_2) \\ & + \bar{s}_1(x_1)s_2(x_2)] [1/4 - (1/3)\sin^2 \theta_w + (2/9)\sin^4 \theta_w] \}, \end{aligned} \quad (3.31)$$

where θ_w is the electroweak angle. Using again the Drell-Yan picture, one finds a similar cross-section to (3.29):

$$\sigma_{Z^0} = (2\sqrt{2}) G_F \pi \tau \int_{\tau}^1 dx z^0(x, \tau/x)/x. \quad (3.32)$$

Before evaluating these cross-sections two comments have to be made concerning the integrals in Eqs. (3.29) and (3.32), called differential luminosity $dL/d\tau$ [49]:

- i) It turns out that the quark distributions measured in deep inelastic lepton scattering [59] are energy-momentum dependent, i.e. scaling is violated for large momentum transfer Q^2 , thus leading to a deviation from the naïve Drell-Yan model. In fact the quark distributions get depleted for large x with increasing Q^2 , whereas they get more populated for small x with increasing Q^2 . This softening of the quark distribution (i.e. structure function) is due to the interaction of quarks, i.e. radiation of gluons leading to additional diagrams [60]. It follows that the quark distribution functions $q(x)$ become $q(x, Q^2)$; in fact they depend logarithmically on the energy scale $Q^2 \approx m_{W,Z}^2$, the energy of the constituent reaction. As a result the quantity $\tau dL/d\tau$ drops for large values of $\sqrt{\tau}$ if this scaling violation is taken into account (Fig. 6a). Note, however, that for the W and Z production ($\tau_w = 0.15$, $\tau_z = 0.17$), there is only a small change between scaling and non-scaling distribution.

In addition, the radiation of hard gluons provides quarks with larger transverse momentum, showing $\langle p_T \rangle \approx \sqrt{Q^2}$. In fact the ISR data [61] show an increase of $\langle p_T \rangle$ as a function of mass. We could expect therefore a rather large value of $\langle p_T \rangle$ for the $W^\pm$ and Z^0 bosons ($\langle p_T \rangle \approx 5$ to 10 GeV/c), which is much larger than expected from the Fermi motion of the constituents of the hadrons.

- ii) It is clear that $p\bar{p}$ collisions offer more constituents for the production of the IVBs since one needs always a quark and an antiquark, the latter being only in the “sea” of the $q\bar{q}$ pairs when

using pp collisions. This is illustrated in Fig. 6b, where $\tau dL/dt$ is roughly a factor of two higher for $p\bar{p}$ collisions than for pp collisions at the relevant value of $\sqrt{\tau}$. The difference becomes more important for large $\sqrt{\tau}$.

With a convenient parametrization of the structure functions including scaling violations and higher-order QCD effects [62] the total cross-sections may be computed [58] (Fig. 7). For $\sqrt{s} = 540$ GeV the values are

$$\begin{aligned}\sigma_{W^\pm} &\approx 1.8 \text{ nb} \\ \sigma_{Z^0} &\approx 0.85 \text{ nb}.\end{aligned}\tag{3.33}$$

However, it has to be pointed out that these values have to be multiplied by a factor k . In fact, it has been observed that the lepton-pair production cross-section is larger than that calculated from distribution functions determined from deep inelastic scattering [56]. This comes from virtual gluon exchange and it has been shown [63] that the second term in the α_s expansion is as large as the first one. The value of k is close to two and constant over the kinematic range $0.02 < \tau < 0.7$, therefore not affecting the shape of the cross-section distributions [64].

3.4 The experimental signatures of the intermediate bosons

As described in Section 3.1, the dominant decay modes of $W^\pm$ and Z^0 are those into $q\bar{q}'$, and $q\bar{q}$, respectively. Pairs of hadronic with large transverse momentum jets should show up. However, QCD backgrounds from quark and gluon scattering are expected to occur 5–10 times more, owing to their strong coupling. Nevertheless, the situation looks more comfortable if one considers heavy-quark jets decaying into light-quark jets. One example is

$$\begin{array}{l} W^- \rightarrow b\bar{t}, \\ \quad \quad \quad \downarrow \\ \quad \quad \quad \begin{array}{l} \rightarrow \bar{b}\ell^-\bar{\nu}_\ell \\ \rightarrow c\ell^-\bar{\nu}_\ell \end{array} \end{array}\tag{3.34}$$

resulting in two leptons of the same sign. Here the only competing background is coming from $b\bar{b}$ and $t\bar{t}$ production. Calculations [65] have shown that the detection of a process such as (3.34) is possible.

The decay of the $W^\pm$ into a charged lepton and its (anti)neutrino, although three times less frequent than the hadronic mode, is far more easy to detect. The appearance of a large transverse momentum lepton, e.g. e or μ , and of a large missing transverse energy as evidence for a neutrino are clear signs for a leptonic decay of a $W^\pm$. In addition, kinematics fortunately conspires in our favour. In fact, the distribution of the transverse momentum p_T of the lepton should show a Jacobian peak (Fig. 8) coming from the production and two-body decay of a resonance. The peak occurs at $p_T = M_W/2$ and is smeared out owing to the finite width of the $W^\pm$, the transverse momentum of the boson, and QCD correction from gluon radiation [66, 67].

Candidates with this kinematical behaviour should show another specific characteristic of the intermediate bosons: a forward–backward asymmetry of the lepton spectrum. Any interference between vector and axial-vector coupling will lead to a definite helicity state of the $W^\pm$ in the production and of the leptons in the subsequent decay (Fig. 9). The weak interaction acts as a

polarizer in the production; the angular distribution of the decay products is a polarization analyser, analogously to the $\pi - \mu$ decay chain. For a $(V - A)$ coupling, inherent in weak interactions, the lepton ℓ^+ will be preferentially emitted in the $\bar{p}$ direction, the ℓ^- in the p direction. It results in a $(1 + \cos \theta^*)^2$ distribution for the charged lepton, where θ^* is its centre-of-mass angle. A distinction between $V - A$ and $V + A$ coupling in production and decay can only be made using polarized p or $\bar{p}$ beams. A helicity flip, i.e. $V - A$ in production and $V + A$ in decay or vice versa, is generally believed to be absent, but can be tested in the $W^\pm$ decays.

The signature for the leptonic decay of the Z^0 is very clean. Two identified leptons with opposite charge and large transverse momentum should give rise to a mass peak. The background, mainly coming from Drell–Yan dilepton production, is expected to be small [66, 67]. As for the $W^\pm$ an asymmetry should also occur in the Z^0 case, but it is rather suppressed owing to the vector coupling to the leptons proportional to $(1 - 4 \sin^2 \theta_w) \approx 0.08$. Therefore only high statistics in Z^0 -production may observe this asymmetry effect.

4. THE $p\bar{p}$ PROJECT

The discovery of the neutral currents [38] in the middle of the seventies supported strongly the unified gauge theories [23, 24], i.e. the existence of three IVBs, W^+ , W^- , and Z^0 . Therefore it was very clear that their experimental observation, as a direct test of the theory, was the way to go. However, the experimentally determined value of $\sin^2 \theta_w$ [68] had led to the expectation that they are very massive, $m_{IVB} > 70 \text{ GeV}/c^2$. Such a high total energy ($\sqrt{s}$) to produce them was not available at that time. Furthermore, the large existing proton accelerators at CERN and FNAL worked as fixed-target machines, i.e. only a small fraction of the beam energy is made available for the creation of new particles. In fact, the total energy

$$\sqrt{s} \approx (2ME)^{1/2}, \quad (4.1)$$

where M is the mass of the target particle ($M \approx 1 \text{ GeV}/c$ in the case of nucleons) and E is the beam energy. As an example, a proton with $E = 300 \text{ GeV}$ (typical for CERN and FNAL) colliding with a proton in the target, gives $\sqrt{s} \approx 25 \text{ GeV}$, a very moderate total energy. A better way to reach higher $\sqrt{s}$ is by the use of a colliding-beam machine, where the accelerated particles meet head-on, giving

$$\sqrt{s} = E_1 + E_2, \quad (4.2)$$

E_1 and E_2 being the momenta of the two beams. The CERN Intersecting Storage Rings (ISR) with $E_1 = E_2 = 32.5 \text{ GeV}$ provided in this way a total energy of 61 GeV , but this is still too low for the creation of the IVBs.

From the above it was quite clear that only colliders would provide the opportunity of producing particles with masses above 50 GeV . Note that for such a value a fixed-target machine would need a beam energy of $E \approx 1250 \text{ GeV}$. Plans for very high energy colliders (i.e. ISABELLE, LEP, etc.) which could produce IVBs were only at a very early stage, their completion being foreseen after 1985.

In 1976 Rubbia et al. [69] came up with an alternative method. Instead of building an entirely new colliding-beam machine, they proposed to convert an existing fixed-target proton accelerator

into a collider. Since such an accelerator has only one annular beam tube it is clear that the counter-rotating beam must be of antiprotons. However, colliders have a disadvantage: the reaction rate is small compared to the one obtained in fixed-target machines. It is $N = L\sigma$ (s^{-1}), where σ is the cross-section and L the luminosity, expressed as follows:

$$L \propto N_p N_{\bar{p}} f / n_b (\epsilon_x \beta_x^* \epsilon_y \beta_y^*)^{1/2} \quad (4.3)$$

where

N_p ($N_{\bar{p}}$) is the total number of (anti)protons/beam

f is the revolution frequency

n_b is the number of bunches (equal for both beams)

ϵ_x (ϵ_y) is the horizontal (vertical) emittance (supposed equal for p and $\bar{p}$)

β_x^* (β_y^*) is the horizontal (vertical) beta value at the interaction point.

To obtain a high luminosity the number of particles per bunch $N_{p,\bar{p}}/n_b$ should be large and both beam dimensions, $(\epsilon_x \beta_x^*)^{1/2}$ and $(\epsilon_y \beta_y^*)^{1/2}$, small. As successfully demonstrated at the ISR, it is very well achieved for protons, leading to $L \approx 10^{31} \text{ cm}^{-2} \text{ s}^{-1}$ with $N_p \approx 10^{12}$.

On the contrary, reaching high $N_{\bar{p}}$ values is not simple. Unlike protons, antiprotons are not readily available from any natural source: they have to be created in high-energy collisions. This happens in an extremely inefficient way. In fact, the ratio of antiprotons (which can be captured into the beam transport system) to protons interacting in a target is

$$N_{\bar{p}}/N_p \approx 10^{-6}. \quad (4.4)$$

For $N_p = 10^{13}/s$ from an accelerator, this gives $\sim 10^7 \bar{p}/s$. By repeating this process over a long period, e.g. a day, values for $N_{\bar{p}}$ of $\sim 10^{11}$ can be achieved. However, all the precious antiprotons have to be stored in the meanwhile before this high $N_{\bar{p}}$ value is reached. In addition, in the production the antiprotons acquire a wide range of directions and momenta. An observer on a moving particle would see the other particles performing randomized motions, such as the heat motion in a gas. Reducing these motions means “cooling the gas”, resulting in a higher particle density in the beam and a smaller size. One speaks of phase-space cooling.

4.1 The beam cooling

The motion of particles is described in a six-dimensional phase space of the three canonical coordinates and the three momenta of the particles. They are subject to Liouville's theorem, which states that for a Hamiltonian system the density of the six-dimensional points, representing the particles in the neighbourhood of any given point, is a constant of motion. In other words, the phase-space volume and the number of particles in the volume are constant. Introducing external forces into this conservative system causes additional degrees of freedom. As an example, radiation emitted by particles of the original system leads to additional coordinates and momenta, and the particles can move separately in incoherent motion. Therefore, phase space can be interchanged between the original system and these new degrees of freedom, i.e. the synchrotron radiation can decrease the phase-space volume of a beam, an effect used in electron accelerators and storage rings.

Two other important non-Liouvillian processes are used today for beam cooling: *electron cooling and stochastic cooling*.

Electron cooling was suggested by Budker already in 1960 [70] for heavy particles such as protons or antiprotons. In this technique, an electron beam travels parallel to a proton beam over a certain distance. The electrons have the same average velocity as the protons throughout the common travelling section. Therefore

$$\langle E_e \rangle = (m_e/m_p) \langle E_p \rangle = (1/1836) \langle E_p \rangle. \quad (4.5)$$

Owing to Coulomb interactions, the protons transfer momenta to the electrons, e.g. they transfer their “heat” to the “cold” electrons. As a result, the proton beam is cooled down, similarly to a gas which becomes cooler when mixed with a colder one. This technique is well understood and has been successfully tested [71].

The principle of *stochastic cooling* was invented in 1968 by van der Meer [72] in order to reduce the emittance of the high-intensity beams circulating in the ISR and thereby increase the luminosity. We give only a general overview and refer the reader to the literature [73, 74]. If information is available about the position in the phase space of a finite number of particles, the particles are no longer completely subject to Liouville’s theorem. In fact, if each particle could be separately observed, its error could be corrected in much the same way as coherent beam instabilities can be corrected by feedback systems [75]. Although it is at present excluded to treat particles independently of each other in beams of 10^5 or more, beam fluctuations caused by the finite number of particles can be viewed as coherent motions and removed by feedback systems, leading to a slow reduction of the phase space.

In practice (see Fig. 10), a non-destructive sensor, the pick-up, measures, for instance, the transverse position of a sample of beam particles. It is connected through a wide-band amplifier to a corrector, the kicker, which corrects this sample. This is done for any successive sample. The sample length is determined by the bandwidth W of the system, which is related to the number N_s of the particles in the sample and the total number N of particles:

$$N_s = N/2WT, \quad (4.6)$$

where T is the revolution time. The system detects and corrects only the average position $\langle x \rangle$ of the particles in the sample. Because of this average only a small amount of cooling can be achieved for each observation, i.e. only a fraction $g\langle x \rangle$ with $g \leq 1$ will be cooled. It can be shown [73, 76] that the cooling rate is

$$1/\tau = (W/2N)(2g - g^2), \quad (4.7)$$

which is largest for $g = 1$. This expression is valid for perfect mixing between kickers (i.e. complete randomizing of the samples within one revolution) and neglecting the amplifier noise. The term $2g$ reflects the coherent effect, the cooling; the term $-g^2$ is the incoherent effect, leading to heating. It is always possible to choose g such that cooling takes place. Non-perfect mixing and amplifier noise affect the incoherent term, so that

$$1/\tau = (W/2N)[2g - g^2(M + U)], \quad (4.8)$$

where $M \geq 1$ is the mixing factor and $U > 0$ represents the amplifier noise. Optimizing Eq. (4.8) leads to an optimum cooling rate

$$1/\tau = (W/2N) [1/(M + U)]. \quad (4.9)$$

This example, applied to transverse (betatron) cooling, gives similar results when considering longitudinal (momentum) cooling, using the momentum dependence of the radial position (Palmer method) or notch filters [77] between a preamplifier and a power amplifier (Fig. 11). In the notch-filter method, the relationship between the momentum of a particle and its revolution frequency yields information about this momentum. A notch filter is a transmission line shorted at its far end, whose length corresponds to half a rotation period. It behaves like a short circuit at all harmonics of the nominal revolution frequency, showing there characteristic “notches”. Particles with a slightly too low momentum are accelerated and those with too high a momentum are decelerated until ideally all particles have “fallen into the notches”.

The stochastic cooling has first been achieved experimentally in the ISR [78] with cooling times of several hours. This successful conclusion led to the construction of the “Initial Cooling Experiment” (ICE) after the suggestion to convert the SPS into a $p\bar{p}$ collider. ICE is a storage ring where both stochastic and electron cooling of proton beams were tested. It confirmed the stochastic cooling theory [73, 79]. Mean cooling times of 15 s for 7×10^7 protons were achieved and the results are compatible with the theory (Fig. 12). Simultaneous cooling in all three planes has yielded lifetimes of about 100 h.

This success in spring 1978 led soon afterwards to the authorization of the $p\bar{p}$ project at CERN. Construction of an intermediate Antiproton Accumulator (AA) using stochastic cooling and the change in operation of the SPS into a $p\bar{p}$ collider with $\sqrt{s} = 540$ GeV were started. A similar project is being performed at FNAL leading to $\sqrt{s} = 2000$ GeV.

4.2 The Antiproton Accumulator

The first step is the production of antiprotons. A beam of 10^{13} protons is accelerated every 2.6 s to an energy of 26 GeV in the Proton Synchrotron (PS) (see Fig. 13) and compressed into one quarter of the PS circumference. This beam is focused on a tungsten target followed by a magnetic horn, collecting 5×10^6 $\bar{p}$ /PS-pulse in the acceptance of the AA. This means that antiprotons will be accepted around a momentum of 3.5 GeV/c, near the production maximum.

The AA ring is a large-aperture fixed-field machine with a vacuum chamber of 70 cm width. This chamber is divided into two parts, one used for the stack and one for the injected beam. Movable ferrite shutters separate these two parts, thus removing the influence of pick-ups and kickers on the two beams. The parameters of the AA are given in Table I.

The cooling and stacking procedure [80, 81] is represented in Fig. 14. The injected $\bar{p}$ pulse is pre-cooled in 2.2 s using the filter method from an initial $\Delta p/p = \pm 7.5\%$ to $\Delta p/p = \pm 0.83\%$. This beam is then trapped in an RF bucket and decelerated to the inner part of the vacuum chamber across the opened shutter, thus building the first part of a stack. The shutter is again closed and a new $\bar{p}$ -pulse is injected after 2.6 s, ready for pre-cooling and subsequent stacking. The stack formed is now continuously cooled both in momentum and in transverse motion, so that particles slowly migrate to the dense bottom of the stack, where the density is four orders of magnitude higher than at the top. In total, 32 pick-ups and about 100 kickers are used for the stack-cooling system. After

24 h of stacking and cooling, 1.6×10^{11} antiprotons have accumulated in the AA. From the high-density part of the stack, on an average 0.6×10^{11} antiprotons [82] are collected in an RF bucket, accelerated to the ejection orbit, sent to the PS for acceleration to 26 GeV/c and transferred to the SPS. The remaining antiprotons in the AA form a new stack and a new accumulation can start.

4.3 The performance of the SPS as a $p\bar{p}$ collider

Once the antiproton stack in the AA is dense enough ($N_{\bar{p}} > 10^{11}$), three proton bunches are injected at 26 GeV/c momentum into the SPS and left there as a coasting beam (see Fig. 13). Then also three bunches of antiprotons, as described above, are injected in the counter-rotating sense into the SPS. Both beams are now accelerated to 270 GeV/c, giving a total energy of $\sqrt{s} = 540$ GeV.

This important achievement was made already three years after the approval of the $p\bar{p}$ project. In fact in December 1981, with a maximum luminosity of $L = 5 \times 10^{27} \text{ cm}^{-2} \text{ s}^{-1}$, the first $p\bar{p}$ collisions were observed [83] and led soon afterwards to the first physical results [84]. Since then important progress has been made in the performance of the involved accelerators (PS, AA, SPS) increasing the luminosity considerably. Table 2 and Fig. 15 give a summary of the performance in 1982 and 1983 [82]. Six experiments, UA1 to UA6 (UA = Underground Area), were accepted [85, 86] to take data at the $p\bar{p}$ collider, but only two of them, UA1 and UA2, are capable of detecting the IVBs. The detectors and results of these two experiments will be described in the next sections.

5. THE EXPERIMENTS UA1 AND UA2

Although a lot of interesting physical phenomena may occur at such a high energy as 540 GeV, the detectors were mainly conceived in view of obtaining the experimental proof of the existence of the IVBs.

From Section 3 the important aspects of such a detector become clear:

- i) It should be as hermetically closed as possible, thus detecting all reaction products.
- ii) It should be capable of detecting the charged leptons: the electrons by their electromagnetic shower cascade and the muons by their penetrating capability through a large amount of material.
- iii) It should visualize all charged tracks and measure their momenta by a magnetic analysis.
- iv) It should measure the energy of the particles by calorimetry. This should provide a detection of the neutrino, identified by the missing energy.
- v) A separation of electromagnetic particles from hadrons should be possible.

All this leads to a general-purpose detector capable of coping with all kinds of physics.

5.1 The UA1 detector

This detector (Fig. 16), covering almost completely the whole solid angle down to a polar angle $\theta = 0.2^\circ$, has the characteristics of track detection and calorimetry. A general description is given in Ref. 87.

5.1.1 The central detector

The centre piece ($5^\circ < \theta < 175^\circ$) is a drift chamber, a cylindrical volume 5.8 m in length and 2.3 m in diameter [88]. This central detector (CD) is located in a dipole magnet generating a field of

0.7 T over a volume of $7 \times 3.5 \times 3.5 \text{ m}^3$. The field direction is horizontal, perpendicular to the beam direction. The CD is a self-supporting cylinder made of six independent half-cylinder chambers. A schematic view of the arrangement of the six modules is shown in Fig. 17 and the most relevant parameters are summarized in Table 3. The orientation of the wire planes has been chosen so as to keep a constant density of points along the tracks over all the detector volume, taking into consideration the expected topology of $p\bar{p}$ interactions and the characteristics of the magnetic field. The drift distance of 18 cm has been chosen, so that with a drift velocity of $5.3 \text{ cm}/\mu\text{s}$ the total drift-time of $3.6 \mu\text{s}$ is less than the time between two $p\bar{p}$ collisions ($3.8 \mu\text{s}$).

The CD has two types of read-out. The drift-time is measured by a 3-bit time-to-digital converter (TDC) interpolator with an accuracy of 4 ns. In addition, the two pulses coming from each wire end are sampled every 32 ns by a 6-bit fast analog-to-digital converter (FADC). This allows a direct measurement of the track position along the wire by the charge-division method. The energy loss dE/dx is measured by the sum of the two pulses through a non-linear scale in order to increase the dynamic range.

The overall spatial resolution is of the order of $290 \mu\text{m}$ (Fig. 18a) in the drift-time coordinate, measured with reconstructed tracks. For the charge-division coordinate one obtains a σ of 1.7% of the wire length (Fig. 18b). This results in a typical relative accuracy of $\pm 20\%$ for a 1 m long track at $p = 40 \text{ GeV}/c$ in the plane normal to the magnetic field. For the ionization loss one derives from the truncated mean of the 60% lowest readings a σ of 6%.

Various calibrations are done before and during the running periods to provide a knowledge of relevant parameters, such as drift velocity, drift angle, absolute reference time t_0 , and absolute wire position. With an electronic test-pulse system different propagation times (e.g. from cables) can be registered to correct the t_0 . With the help of X-rays and ionizing laser beams, simulating tracks, the drift velocity and the drift angle can be precisely measured. As a last method one uses real tracks, e.g. cosmic-ray muons. Since those tracks are normally crossing different gaps and also different modules, it is possible to fix the parameters by making a best fit to the tracks.

5.1.2 The electromagnetic calorimeter

Next to the CD is the electromagnetic calorimeter, still inside the magnet. It consists of two parts (see Table 4):

- i) A cylindrical calorimeter, 6 m long with an inner diameter of 2.72 m, covers the region $25^\circ < \theta < 155^\circ$ [89]. It consists of 48 semicylindrical half shells, called “gondolas”, 24 on each side of the beam axis. The width of each gondola is 22.5 cm (Fig. 19). Alternating layers of scintillator (1.5 mm) and lead (1.2 mm) form four segments in depth of a gondola. The radiation length for the four segments is the following: $3.3/6.6/9.9/6.6 X_0$, resulting in a total of $26.4 X_0$ and 1.1 interaction lengths, thus scarcely affecting the hadrons. The scintillation light at 380 nm from all four segments is absorbed by wavelength shifters, re-emitted at 480 nm and transmitted to photomultipliers via light-guides. Each segment is seen by four photomultipliers located outside the magnet. Two amplifiers are connected to each anode. The first one has two outputs, one for low gain and one for high gain. This is necessary owing to the wide momentum spectrum of the particles. The second amplifier integrates the pulses and is used in the trigger logic. One benefits from the light attenuation in order to improve the localization of the deposited energy. Pairing the top and bottom photomultipliers gives an azimuthal angle ϕ with a resolution $\sigma(\phi) = 0.3 \text{ E}^{-1}$ (Fig. 20a). Similarly, pairing left and right photomultipliers along the beam

direction determines this coordinate (x) to $\sigma(x) = 6.3 \text{ E}^{-1}$ (Fig. 20b). These results were obtained using a test beam of electrons [89]. In addition, it served to calibrate the energy of the gondolas and determine the energy resolution to $\sigma(E) = 0.15 \text{ E}^{-1}$ (Fig. 21). Also the separation between electrons and pions was studied by observing the deposited energy in the hadron-calorimeter module placed behind the gondola. The result is shown in Fig. 22 for normal incidence of the test beam. Once the detector was installed, energy calibrations were performed periodically with a strong ^{60}Co source (7 Ci), by making detailed scans over the whole electromagnetic detector. The stability of the gain of the photomultipliers was monitored using a laser system, which in turn is monitored by a ^{241}Am source and photodiodes.

- ii) The end faces of the CD are covered by an end-cap electromagnetic calorimeter (“bouchons”) in a region $5^\circ < \theta < 25^\circ$ (Fig. 23). Each “bouchon” is divided into two parts, each part being divided azimuthally into 16 equal sectors (petals). Each petal, a sandwich of lead (4 mm) and scintillator (6 mm), is segmented into four parts in depth with 3.6, 7.2, 8.7, and 7.2 radiation lengths. As for the gondolas, wavelength shifters transmit the light to the photomultipliers. Each segment has one photomultiplier. The attenuation length of the scintillator has been chosen to match the variation of $\sin \theta$ over the radius of the calorimeters, so as to measure directly in first approximation $E_T = E \sin \theta$ rather than the true deposition E . This permits the amount of transverse energy deposited to be read out directly. After the first two segments a position detector is located to measure the shower position. It consists of two planes of orthogonal proportional tubes of $2 \times 2 \text{ m}^2$ cross-section; it locates the centre of gravity of electromagnetic showers to $\pm 2 \text{ mm}$ in space and measures the energy deposited.

5.1.3 The hadron calorimeter

The hadron calorimeter (Fig. 19) covers the angular range $5^\circ < \theta < 175^\circ$ and is formed of two types of modules: in the central region one finds 16 C-shaped modules, 8 on each side of the beam; in the front region 12 I-shaped modules close the detector. The C’s are also the return yoke of the magnet. The whole calorimeter (Table 4) is of the iron/scintillator sandwich type: 1 cm scintillator plates alternate with 5 cm iron plates. Each C is subdivided into 12 sections in azimuth and two equal parts in depth: a front stack with 7 (6 in the corners) scintillator plates separated by the iron plates and a rear stack with 8 scintillator plates. As for the electromagnetic calorimeter, the light is transferred via wavelength shifters and light-guides to the photomultipliers, two per stack. The I’s are subdivided into 2×6 stacks but the 16 stacks nearest the beam are further subdivided into four. The number of scintillators is now 12 for front stacks and 11 for rear stacks, all separated by 5 cm iron plates. The total thickness of iron is therefore greater than for the C’s because of the larger average energy in the forward region.

Each photomultiplier signal is digitized using 12-bit ADCs. For triggering on hadronic energy the 12 photomultiplier signals from a 45° section of a C are summed, digitized to 8-bit accuracy, and then weighted by the average value of $\sin \theta$ for these stacks to give transverse energy. In addition, the two photomultiplier signals from each rear stack are added and used in the muon trigger to select minimum ionizing particles.

The calorimeter is monitored by a fibre-optics system which distributes ultraviolet (UV) light from a nitrogen laser to each scintillator plate. A test module has been extensively studied in a test beam using hadrons and muons [90]. There the resolution was found to be $\sigma = 0.8^{-1}$ (Fig. 24).

The total thickness of the calorimeter, 80 cm iron for the C's (125 cm for the I's) and the lead of the gondolas, is 6.1 absorption lengths, sufficient to contain almost totally a hadronic shower.

5.1.4 The muon detector

In order to complement the electron identification over as large a solid angle as possible a muon detector [91] covers the whole UA1 experiment. The muons are identified in a large drift-chamber array, after penetration of the $6.1/\sin \theta$ absorption lengths of the calorimeters. An additional iron shielding has been placed on the top (40 cm) and the side (60 cm) of the magnet, in front of the muon detector. The absence of secondary interaction inside the absorber is signalled (i) by the exiting coordinates and angles being compatible with what is expected from magnetic deflection and multiple scattering and (ii) by the minimum ionization of the track in the calorimeter segments traversed.

Fifty muon chambers, nearly $4\text{ m} \times 6\text{ m}$, surround the whole detector. Each chamber consists of two orthogonal layers of drift tubes with two places per projection (Fig. 25). With a staggered arrangement for adjacent planes one resolves the left-right ambiguity and at the same time compensates for the inefficiency from the dead spaces between tubes. The extruded aluminium drift tubes have a cross-section of $45\text{ mm} \times 150\text{ mm}$, leading to a maximum drift length of 70 mm. The main parameters of the muon detector are summarized in Table 5. The sense wires are connected to a charge-sensitive amplifier followed by a voltage-sensitive amplifier [92]. The output signals are sent over twisted pairs to multitime digitizers (MTD), which convert the analog pulses into digitized pulses with a precision of 8 ns [93]. Tests in a beam and with cosmic-ray muons have shown that an average spatial resolution of $\sigma = 300\text{ }\mu\text{m}$ is achieved throughout the sensitive volume of the tubes.

Owing to the high interaction rate at the $p\bar{p}$ Collider, the muon detector itself must provide a trigger on high-momentum muons in the presence of a severe background. Therefore, two chambers, of four planes each, have been placed on five sides of the detector at a distance of about 60 cm (Fig. 25). This long lever arm leads to an angular resolution of a few milliradians, comparable to the average multiple-scattering angle of high-energy muons, typically 3 mrad at 40 GeV/c. The remaining side, the bottom, was closed with special chambers consisting of four parallel layers of drift tubes. This whole muon detector, covering an area of about 670 m^2 , permits an easy definition of a track pointing to the interaction vertex within a specified cone of aperture $\pm 150\text{ mrad}$. This is accomplished by a dedicated set of hardware processors, the muon fast trigger, filtering the pattern of the tube hits. The operation is carried out in less than $1\text{ }\mu\text{s}$ after the maximum drift-time in the muon chambers ($1.5\text{ }\mu\text{s}$) has elapsed and it introduces no additional dead-time, since it is accomplished before the next beam crossing. The forward chambers have been only partially used in the trigger, since their trigger rate was too high.

5.1.5 The forward detectors

Additional calorimetry, both electromagnetic and hadronic, and track detection extend to the forward regions of the experiment, down to 0.2° (see Fig. 16b and Table 4).

First, drift image chambers, the "end-cap chambers", are located in the I's of the hadron calorimeter. They are of trapezoidal cross-section, one chamber above the beam pipe and one below, and this for both sides of the experiment. They have two sense-wire planes with drift lengths of 19 cm and 9 cm. Drift-time measurement and charge division give the spatial coordinates of the tracks. They cover the two polar angular regions $1^\circ < \theta < 5^\circ$ and $175^\circ < \theta < 179^\circ$.

Outside the UA1 detector one finds on each side a calorimetrized compensator magnet (“CALCOM”). The electromagnetic calorimeter is subdivided into four segments in depth, each one made out of 13 lead plates (3 mm) and 14 scintillator plates (3 mm) (see Table 4). After the first segment, a position detector (made of proportional chambers) registers the electromagnetic shower. As for the C’s and I’s, the return yoke of the compensating magnet is part of the hadronic calorimeter. Its 4 cm iron plates are followed by 8 mm scintillator plates. The whole calorimeter is segmented in 6 parts in depth with a position detector after the first one.

The “very-forward” detector is made of a drift chamber, followed by an electromagnetic and then a hadronic calorimeter (Table 4). (For more details see Ref. 87.)

Finally at ~ 25 m from the intersection point, drift chambers (“Roman pots”) are installed. They can be positioned to within a few millimetres from the beam, covering therefore angular ranges of $0.01^\circ < \theta < 0.1^\circ$ and $179.9^\circ < \theta < 179.99^\circ$. They are used for the measurement of elastic $p\bar{p}$ scattering [94] and thus serve as a luminosity monitor.

5.2 The UA2 detector

This detector (Fig. 26) is instrumented to identify electrons and measure their energy over a polar angle interval $20^\circ < \theta < 160^\circ$, but it does not detect muons. Therefore the central region has no magnetic field; only for $20^\circ < \theta < 37.5^\circ$ and $142.5^\circ < \theta < 160^\circ$ are there magnetic spectrometers for the purpose of measuring the important charge asymmetry in the $W \rightarrow e\nu$ decay.

5.2.1 The vertex detector

This detector (Fig. 27), surrounding the $p\bar{p}$ interaction point, is a system of cylindrical chambers which measures charged-particle trajectories in a region without magnetic field. It consists of two types of chambers:

- i) Four multiwire proportional chambers, VS1 to VS4, with digital read-out on 2112 anode wires and analog read-out on 2160 cathode strips, at $\pm 45^\circ$ with respect to the wires.
- ii) Two large-gap drift chambers, VJ1 and VJ2, with 12 sense-wire planes measuring the track position by charge division to 1% of the wire length.

These chambers are used to determine the position of the event vertex along the beam line with a precision of ± 1 mm and the ionization loss to 20%.

The vertex detector is surrounded by a 1.5 radiation lengths tungsten converter followed by a proportional chamber C5, the pre-shower counter. It is equipped for charge division read-out on 288 wires and pulse-height measurement on 960 helical strips leading to an accuracy in position determination of about 10 cm along the beam and a track-matching accuracy of $\sim \pm 3$ mm. This counter provides the desired accuracy in angular measurement and discriminates against events in which a charged hadron and one or more photons hit the same calorimeter cell.

5.2.2 The central calorimeter

This calorimeter contains an electromagnetic and a hadronic part and covers the full azimuth and a polar angle interval $40^\circ < \theta < 140^\circ$ [95, 96]. A segmentation of 240 cells is chosen. Each cell (Fig. 28) covers 10° in polar angle and 15° in azimuth and has its front face located 60 cm away from the centre of the $p\bar{p}$ interaction region. The cell has the shape of a truncated pyramid with its (imaginary) peak in the centre of the interaction region in order to minimize the energy sharing between adjacent cells. Ten of those cells form a mechanically rigid unit of orange-slice shape (Fig. 26).

Each cell is segmented in three longitudinal sections (Fig. 28):

- i) The electromagnetic compartment, 17.5 radiation lengths deep, consists of 26 lead plates (3.5 mm thick) and 27 scintillator plates (4 mm thick).
- ii) Two hadronic compartments, each roughly 2 absorption lengths deep, use 19 iron plates (15 mm thick) in the first compartment (21 in the second) interspaced by 5 mm thick plates of acrylic scintillator.

The total cell length corresponds to ~ 4.5 absorption lengths, barely sufficient to completely contain high-energy hadronic showers with energies larger than 40 GeV/c.

The light is collected for all three compartments by wavelength shifters on both sides of the calorimeter cell, and transmitted via light-guides to photomultipliers. The analog signals from the photomultipliers are converted into digital form by charge-sensitive 12-bit ADCs.

Each cell of the central calorimeter has been calibrated in a 10.10 GeV/c beam using electrons for the electromagnetic calorimeter and muons for the hadronic one. The calibration stability has been measured continuously by a light-flash system, using a xenon flash tube, and a movable 4 mCi ^{60}Co radioactive source, running along the front faces of the calorimeter. The systematic uncertainty in the absolute energy scale is less than $\pm 1.5\%$ for the electromagnetic and less than 3% for the hadronic calorimeters, with channel-to-channel r.m.s. spreads of 3% and 8%, respectively.

The response of both calorimeters with respect to beam impact position, beam incidence angle, and beam energy, has been studied using electron and pion beams of momenta between 1 and 70 GeV/c. For electrons the corrected energy is

$$E_{\text{corr}} = 1.024 E_{\text{meas}} / [1 + 0.01 \ln(E_{\text{meas}} + 1)] \quad (5.1)$$

(E in GeV) and can be measured to better than $\pm 1\%$ on an average. The energy resolution (Fig. 29) is measured to be $\sigma(E) = 0.14 E^{1/2}$. From the longitudinal shower-development study, it turned out that, at 70 GeV electron energy, 95% of all electrons deposit less than 11% of their energy in the first hadron calorimeter compartment (Fig. 30). Electrons of normal incidence at the centre of the cell lose less than 1.5% of their energy into adjacent cells owing to the lateral shower development. For electrons entering the edge of the cell, it is possible, knowing the track impact position, to predict the energy fractions deposited in the hit cell and the surrounding ones.

For the hadron calorimeter, data collected with incident π^- beams have been used to determine the ratio r between E_{meas} and E . The dependence on energy of the mean value of r , $\bar{r}$, and its r.m.s. deviation from $\bar{r}$, $\text{rms}(r)$ are shown in Fig. 31a,b. The resolution $\text{rms}(r)/\bar{r} (\triangleq \sigma/E)$ decreases from $\sim 30\%$ at 1 GeV to $\sim 13\%$ at 70 GeV (11% for normal incidence at centre of cell), approximately proportional to $E^{-1/4}$.

5.2.3 The forward detector

The two forward regions (see Fig. 26) are each equipped with twelve toroidal magnet sectors having a mean-field integral of $0.38 \text{ T} \cdot \text{m}$. Each sector is instrumented with:

- i) Three drift chambers (Fig. 32) located after the magnet [97]. Each chamber has three sense planes with wires oriented at -7° , 0° , and $+7^\circ$ with respect to the magnetic-field direction. The point resolution on tracks is $\sigma_p = 200 \mu\text{m}$ and the angular resolution in the bending plane is $\sigma_t = 0.47 \text{ mrad}$. When combined with the vertex detector one gets an accuracy in momentum

determination of $\sigma \leq 0.02 \times p^2$ (p in GeV/c). The total system of 72 chambers has 216 independent planes with 2280 sense wires.

- ii) A 1.4 radiation length thick lead-iron converter, followed by a multitube proportional chamber (MTPC) (see Fig. 32). The trapezoidal-shaped chamber [98] consists of three layers (wire orientation $-7^\circ, 0^\circ, +7^\circ$) of drift tubes with two planes per projection. Adjacent planes have a staggered arrangement. The cylindrical brass tubes have an inner diameter of 19.4 mm and 226 of them form one chamber. Electromagnetic showers are localized with a precision of ≤ 8 mm.
- iii) An electromagnetic calorimeter of the lead-scintillator sandwich type, subdivided into 10 independent cells. Each cell has 2 segments in depth, 24 and 6 radiation lengths thickness. The leakage for hadrons into the second segment is $< 1\%$. Each segment is read out by two wavelength shifters. The energy resolution is $\sigma \approx 0.15 E^\pm$ and the calibration stability is better than 3%.

6. THE DATA

We report on the data taken at the end of 1982 and in spring 1983. During these two Collider periods, the initial luminosity increased from 3×10^{28} to $1.5 \times 10^{29} \text{ cm}^{-2} \text{ s}^{-1}$ and the two beams were left on the average for 12 h in the Collider before a new refill was done. A total of about 8×10^9 collisions at $\sqrt{s} = 540 \text{ GeV}$ took place in this time.

6.1 The UA1 experiment

This experiment took data for an integrated luminosity of 136 nb^{-1} for the $W \rightarrow e\nu$ and $Z^0 \rightarrow \ell^+ \ell^-$ ($\ell = e, \mu$) decays and 108 nb^{-1} for the $W \rightarrow \mu\nu$ decay. These values are obtained after subtraction of dead-time and other instrumental inefficiencies and are known to about $\pm 15\%$ uncertainty.

6.1.1 The trigger and the data taking

A pre-trigger from a coincidence of scintillation hodoscopes in the forward arms of the detector selected the $p\bar{p}$ interactions and rejected by timing the beam-gas interactions. This pre-trigger was used either for minimum-bias data taking or as the start of the main-trigger system. Here four types [43, 99] were operated simultaneously.

- i) An “electron trigger”, namely at least 12 GeV of transverse energy deposited in two adjacent elements of the electromagnetic calorimeters covering angles larger than 5° with respect to the beam line.
- ii) A “muon trigger”, namely at least one penetrating track detected in the muon chambers with pseudorapidity $|\eta| \leq 1.3$ and pointing in both projections to the interaction vertex within a specified cone of aperture $\pm 150 \text{ mrad}$.
- iii) A “jet trigger”, namely at least 20 GeV of transverse energy in a localized calorimeter cluster. A cluster is defined either as a group of eight gondolas and the two hadron calorimeter elements immediately behind, or as a quadrant of bouchon elements with the corresponding hadron calorimeters behind them.
- iv) A global “ E_T trigger”, with more than 50 GeV of total transverse energy from all calorimeters with $|\eta| < 1.4$.

These triggers, except the “muon trigger”, were built up by two identical trigger processors, transforming pulses to energies and making the relevant sums. The “muon trigger” came from the muon fast trigger. Once an event is accepted by the processors, the read-out, using the CERN Romulus Remus system, starts to transfer all the information from approximately 200 CAMAC modules to the control room. This is done by reading in parallel five branches containing the whole of the data. Before the data goes on to tape at the main computer, a Norsk Data ND 100/500, four 168E computers select, by the so-called express line, events with an “electron trigger” flag, dimuons and cosmic-ray muons (for calibration). These events are written on a special tape. At different places of the read-out it is possible to spy into the data transfer by means of small computers for monitoring reasons. More information about the complete read-out can be found elsewhere [87, 89, 100]. The event rate with the four types of triggers was about 3 Hz at a luminosity of $10^{29} \text{ cm}^{-2} \text{ s}^{-1}$.

6.1.2 Lepton identification and event selection

The events of the express line have been fully reconstructed for $W^\pm$ and Z^0 search. At first we consider the electrons. In order for an electromagnetic cluster in the calorimeters to be accepted as an electron, there has to be a CD track that points to it, i.e. the impact point of the extrapolated track goes with the centroid of the localized electromagnetic energy deposition. In addition, isolated electrons are identified by their characteristic transition curve, and in particular by the lack of penetration in the hadron calorimeter (see Section 5.1.2 and Fig. 22). The misinterpretation of an electron as a pion is less than 1% at 40 GeV/c. For the $W^\pm$ search also neutrino identification plays an essential role. The apparent visible energy imbalance (missing energy) is indicative of the existence of a neutrino. Here the calorimetry down to 0.2° in polar angle θ is of fundamental importance. One can define an energy flow vector $\Delta\vec{E}$, adding vectorially the observed energy depositions over the whole solid angle (including muons). Neglecting particle masses, and with an ideal calorimeter response and solid-angle coverage, momentum conservation requires $\Delta\vec{E} = 0$. For minimum-bias and jet-enriched events for which neutrino emission ordinarily does not occur, the transverse components ΔE_y and ΔE_z exhibit small residuals centred on zero, with an r.m.s. deviation described by $\Delta E_{y,z} = 0.4 \sqrt{E_T}$, where E_T (in GeV) is the scalar sum of all transverse energy contributions in the event. The distributions have Gaussian shape and no prominent tails (Fig. 33). The longitudinal component is of little use, since it is strongly affected by energy flow escaping undetected through the beam pipes.

After the complete off-line reconstruction one applies technical cuts according to the above criteria. One requires three conditions:

- i) A fast track with $p_T \geq 7 \text{ GeV/c}$ (or compatible, within 3σ , with $p_T \geq 15 \text{ GeV/c}$) recorded by the CD must hit an electromagnetic cluster, defined by a pair of gondolas or bouchon petals. The position information from the e.m. calorimeter must agree with the measured impact point of the charged track. In the case of the gondolas, the azimuthal measurement must agree within 5σ ; in the case of the bouchons the following inequality must be verified:

$$(y_{CD} - y_B)^2/25 + (z_{CD} - z_B)^2/225 < 1, \quad (6.1)$$

where y_{CD} , z_{CD} and y_B , z_B are the position (in mm) measured in the CD and in the position detectors of the bouchons, respectively.

- ii) This track must be isolated. Additional particles in the neighbourhood might fake the energy measurement of the electron. Therefore, other tracks entering the same pair of gondolas or bouchon petals must not add up to more than 2.5 GeV/c of transverse momentum. In addition, the track should not belong to a jet. For this reason, additional charged particles lying within a cone $\Delta r = (\Delta\eta^2 + \Delta\phi^2)^{1/2} < 0.7$ around the fast track, may contribute at most 10% of the transverse energy contained in the e.m. cluster.
- iii) The shower profile in the two gondolas or two petals must be consistent with that of an electron, as measured in the test beam. The shower should be entirely contained in the e.m. calorimeter, i.e. the energy deposition E_{had} in the hadronic calorimeter immediately behind should not exceed 600 MeV in the case of the gondolas and 1500 MeV in the case of the petals. In addition one requires that the energy in the fourth segment be less than that in the third segment, i.e.

$$E_4 < 500 \text{ MeV} \quad \text{or} \quad E_4/(E_3 + E_4) < 0.2. \quad (6.2)$$

More than 98% of all electrons satisfy this requirement.

From the initial 150,000 reconstructed events with an electron trigger, 346 remained [99] after these cuts and were individually examined by physicists on the visual scanning and interactive facility Megatek.

As described in Section 5.1.4, the muon proves to be a non-interacting particle leaving the detector unaffected (apart from multiple scattering). To be accepted as a muon candidate, the extrapolated CD track must agree in position and angle with the one measured in the muon chamber (Fig. 34), whereby the distance of ~ 6 m to the vertex helps considerably. This extrapolation has been extensively studied with high-momentum ($p > 10$ GeV/c) cosmic-ray muons. The track fitted in the CD and its errors $\Delta(1/p)$, $\Delta\lambda$, and $\Delta\phi$ (Fig. 34) were extrapolated through the calorimeters and the additional iron up to the muon chambers. To these extrapolated CD errors, the errors on the track reconstructed in the muon chamber, and terms taking into account the multiple scattering in angle and displacement were added in quadrature. A χ^2 value was calculated for the deviation both in position and angle, adding a constant error (5 cm in position and 20 mrad in angle) to take into account unknown systematics. A high- p_T track was accepted by the selection, if the matching satisfied a cut of $\chi_a^2 > 25$ in angle and $\chi_d^2 > 25$ in position.

The track measurement in the muon chambers permits a second measurement of the particle momentum because of two effects: the muon is again deflected in the magnetic field of the flux return yoke and the track measurement in the CD is improved owing to the long lever arm with the muon chambers. The statistical and systematic errors in this second measurement were carefully checked with high-momentum cosmic-ray muons. Figure 35 shows a comparison of the momenta as measured in the CD alone and with the muon chambers in addition.

For the neutrino identification the same arguments are true as for the electron case, except that now the transverse momentum of the muon has to be subtracted vectorially from the missing-energy vector in order to obtain the transverse momentum of the neutrino.

Since single muons (and multimueons) are only on the original tapes one is faced with 2.5×10^6 events. Dimueons are on the express-line tape but we will follow the selection from the original tapes since there also the dimueons are present. A special filter program was set-up to select muons. Only 10^6 events had a “muon trigger” flag. Those events are contaminated by hadron leakage through the absorber, beam halo, meson decays, and cosmic rays. The program reconstructed

tracks in the muon chambers. If they pointed roughly towards the interaction vertex, the CD information was decoded along a path from the muon chamber hits to the vertex region. Track finding and fitting were performed in this path. Events were kept if the CD track satisfied a momentum cut of $p_T > 3 \text{ GeV/c}$ or $p > 6 \text{ GeV/c}$ and matched the muon track within generous limits. After this filter, 72000 events were left. Only these events with $p_T > 5 \text{ GeV/c}$ were then completely reconstructed, leading to 17000 events. From these events, 689 had at least one muon candidate with $p_T > 15 \text{ GeV/c}$ or $p > 30 \text{ GeV/c}$ [101]. They were examined on the Megatek visual display.

6.2 The UA2 experiment

The data sample corresponds to an integrated luminosity of 131 nb^{-1} . Only the search for electrons can be done.

6.2.1 The trigger

Three different types of trigger [44, 46, 102, 103] were operated simultaneously:

- i) A W-trigger, namely at least 8 GeV of transverse energy deposited in matrices of 2×2 cells of the central electromagnetic calorimeter. In the forward electromagnetic calorimeters, only those matrices made up of cells belonging to the same sector were considered.
- ii) A Z^0 -trigger, namely at least 3.5 GeV of transverse energy deposited in two matrices at a time and separated in azimuth by more than 60° .
- iii) A “jet trigger”, namely at least 25 GeV (later 35 GeV) of transverse energy (electromagnetic and hadronic) in coincidence with a minimum-bias trigger.

6.2.2 Electron identification and event selection

The presence of an electron is signalled, as in the UA1 experiment, by an electromagnetic shower development consistent with an electron. In particular, electrons are identified by a small lateral size in the electromagnetic compartment of the calorimeter and limited leakage in the hadronic compartment. This is satisfied by more than 95% of isolated electrons, as verified experimentally with test-beam data. Neutrinos identify themselves by the presence of a large missing transverse energy.

The full data sample consists of $\sim 7 \times 10^5$ W triggers and $\sim 7 \times 10^5$ Z^0 triggers. A first data reduction is done for both classes by requiring:

- i) for the $W^\pm$, the presence of an energy cluster with a transverse electromagnetic energy greater than 15 GeV;
- ii) for the Z^0 , a total electromagnetic transverse energy exceeding 30 GeV and the presence of a pair of energy clusters (electromagnetic and hadronic) having an invariant mass in excess of 50 GeV/c.

In the central calorimeter, clusters are obtained by joining all electromagnetic cells which share a side and contain at least 0.5 GeV, including also those cells having at least one side in common with a cluster. In the forward calorimeter a cluster consists of at most two adjacent cells having the same azimuth. The cluster energy is the sum of deposited electromagnetic and hadronic energy.

The invariant mass of a cluster pair in the Z^0 -sample is calculated under the assumptions that the event vertex is in the centre of the detector and that the particles impinge on the cluster centroid, thus determining their momenta.

With this preselection 8.3×10^4 events remain in the W-sample and 7427 events in the Z^0 -sample. All these events are then fully reconstructed.

These samples are now subject to additional cuts which filter out the electron candidates. They are summarized in Table 6. The cut efficiencies η are evaluated using the measurements of the response of the various parts of the detector in the electron test beam [96–98]. The overall efficiency, i.e. the product of the individual efficiencies η , is $(76 \pm 5)\%$ for the central region and $(80 \pm 5)\%$ for the forward regions. Cut 1 reduces the W-sample to 10169 events (8192 in the central region and 1977 in the forward regions) and the Z^0 -sample to 24 events. Applying the other cuts reduces the W-sample further to 225 events with 227 electron candidates (200 and 27 for the two parts of the detector).

7. SEARCH FOR THE W^+ AND W^-

From this initial event selection it remains to crystallize the $W^\pm$ -candidates, observing their decays $W \rightarrow e\nu$ and $W \rightarrow \mu\nu$, and compare them with expectation.

7.1 The decay $W \rightarrow e\nu$

Both experiments, UA1 and UA2, have searched for this decay channel and found evidence for it at the beginning of 1983 [42–44].

7.1.1 The UA1 experiment

The remaining 346 events [99, 104–107] with isolated electrons (see Section 6.1.2) have been classified after scanning into two categories:

- i) 291 events show a clear jet [108, 109] opposite to the electron track within 30° in azimuth. These events are highly contaminated by jet–jet events where one jet fakes the electron signature. They are rejected. Besides they have essentially no missing energy.
- ii) 50 events have no jet at all and 5 events have no jet opposite to the electron within 30° . These 55 events are characterized by a significant missing energy (Fig. 36), indicating the presence of neutrino emission. According to the experimental energy resolutions, at most the three lowest missing energy events are compatible with no neutrino emission. They are excluded by applying a cut on the missing energy of 15 GeV. On the other hand, an alternative analysis based exclusively on the presence of missing transverse energy yields the same 52 events.

These events show a clear electron signature. In fact, the energy deposited in the hadronic calorimeter behind the gondolas has values far away from the imposed cuts (Fig. 37c) and the fraction of electromagnetic energy deposited in the first and fourth gondola segments (Fig. 37a,b) fall within expectations from the test-beam measurements. In addition, the geometrical matching between the CD track and the cluster in the gondolas is quite perfect (Fig. 38), thus supporting the absence of composite overlaps of a charged track and π^0 's expected from jets. In order to ensure the best accuracy in the electron energy determination, only events in which the electron track hits the gondolas more than $\pm 15^\circ$ away from their top and bottom edges have been retained. This reduces the sample to 43 events, for which the momentum measured by the CD agrees, within errors, with the energy deposited in gondolas (Fig. 39). A typical $W \rightarrow e\nu$ candidate is shown in Fig. 40, where the raw digitizings of the CD are represented. Note the clear, isolated, well-measured straight track of the electron (Fig. 40a) and the clean electromagnetic energy of it in the absence of any hadronic energy (Fig. 40b).

Before discussing more physical properties of these 43 W-candidates, we estimate the possible sources of background [89, 99, 104, 106]:

- i) Misidentification of a high- p_T charged pion as an electron or overlapping of one or more π^0 's with a high- p_T charged pion. In the full data sample the UA1 group has searched for isolated high- p_T charged hadrons satisfying the same isolation criteria as the electron candidates. However, this time no cuts are imposed on the shower profile in the gondolas, but $E_{\text{had}} > 600$ MeV was required in order to ensure hadronic signature. After imposing no jet activity, there remain 8 events with $E_T > 15$ GeV. Using the pion rejection values from the test-beam measurement (see Section 5.1.2 and Fig. 22), the resulting background amounts to < 0.5 events. However, some background events may be detected from other decays of the W, namely

$$\begin{aligned} W &\rightarrow \tau \nu_\tau \\ &\quad \hookrightarrow \pi^\pm (\pi^0) \nu_\tau \end{aligned} \quad (7.1)$$

or

$$\begin{aligned} W &\rightarrow \tau \nu_\tau \\ &\quad \hookrightarrow e \nu_e \nu_\tau \end{aligned} \quad (7.2)$$

amounting to < 2.5 events, but contributing only to the low- p_T part of the electron spectrum. An additional test proves the negligibility of the pion background. One can estimate the inclusive hadron spectrum using the measured high- p_T π^0 spectrum in the region $1.5 < |\eta| < 3.0$. Assuming $\pi^0 = \frac{1}{2}(\pi^+ + \pi^-)$, this yields a background of 0.3 events.

- ii) High- p_T π^0 , η^0 , or γ converting internally (Dalitz decay) or externally to an e^+e^- pair with one track missing.

The number of isolated electromagnetic conversions (π^0 , η , γ , etc.) per unit of rapidity has been directly measured as a function of E_T in the “bouchons” in the interval 10–40 GeV. From this spectrum, the Bethe–Heitler formula for pair creation, and the Kroll–Wada formula for Dalitz pairs [110], the expected number of events with a single $e^\pm$ with $p_T > 20$ GeV/c is $1.4 p_0$ (in GeV/c), where p_0 is the effective momentum below which the low-energy track of the pair becomes undetectable. With $p_0 = 200$ MeV/c (very conservative), this background is negligible.

Having demonstrated that there is no background we conclude that the 43 events have indeed high- p_T electrons (Fig. 41). This distribution shows a clear Jacobian peak as expected from a $W \rightarrow e\nu$ decay (see Section 3.4). Furthermore, the large missing energy (Fig. 42) observed in all events, interpreted as being due to the emission of one (or several) non-interacting neutrino(s), exhibits the same Jacobian peak. If the assumption of a $W \rightarrow e\nu$ decay is correct, there must be a strong correlation in angle and energy between electron and neutrino. In fact, this is clearly observed. In Fig. 43 the vector momenta of the electron and the neutrino are shown to be almost exactly balanced back-to-back. In other words, the component of the neutrino energy parallel to the electron direction shows the same magnitude as the electron transverse energy (Fig. 44a) and the neutrino energy transverse to the electron direction is small and independent of the electron transverse energy (Fig. 44b). This behaviour is characteristic for the decay of a massive, slow particle, whose mass is approximately twice the maximum observed transverse momentum of decay products.

Note, however, that the high- p_T end of these spectra are affected by the transverse momentum of the W. This transverse momentum is simply obtained by adding the transverse momenta of the lepton pair, $\vec{p}_T^W = \vec{p}_T^e + \vec{p}_T^{\nu}$. Its distribution is shown in Fig. 45, leading to an average value $\langle \vec{p}_T^W \rangle = 6.3 \text{ GeV}/c$. The five events, which have a visible jet, have also the highest value of $\vec{p}_T^W$. But this is approximately balanced by the p_T of the recoiling jet. These events are indeed expected because of the additional higher-order QCD diagrams (Section 3.3) involving gluons [67]. The experimental distribution is in good agreement with the many theoretical expectations from QCD for the production of a massive state via the Drell-Yan quark-antiquark annihilation [62, 111–113].

For the determination of the W-mass m_W three methods, using maximum likelihood fits, were applied:

- i) It can be obtained from the transverse-momentum distribution of the electrons (Fig. 42). Note, however, that this variable can be appreciably affected by the transverse momentum of the W. One obtains $m_W = (80.5 \pm 0.5) \text{ GeV}/c^2$, taking QCD predictions into account.
- ii) One can define a transverse mass of the W-particle, $m_T^W = [2E_T^e E_T^{\nu} (1 - \cos \phi_{e\nu})]^{1/2}$, where $\phi_{e\nu}$ is the angle between E_T^e and E_T^{ν} (Fig. 46a). For a zero longitudinal component of the W, the variable m_T^W is equal to m_W ; in all other cases it is smaller. The advantage of this method is that m_W is almost independent of the transverse momentum of the W and the angular distribution of the charged lepton. It may be, however, that the lower part of the m_T^W distribution is affected by $W \rightarrow \tau \nu_\tau$ decays. The result from the likelihood fit gives $m_W = (80.3_{-1.3}^{+0.4}) \text{ GeV}/c^2$.
- iii) One can overcome the small disadvantage mentioned above by imposing that E_T^e and E_T^{ν} are larger than $30 \text{ GeV}/c$. The resultant distribution (Fig. 46b) then shows a relatively narrow peak, at approximately $76 \text{ GeV}/c^2$. The fit yields now $m_W = 80.9 \pm 1.5 \text{ GeV}/c^2$. This value should be retained as the best one, since it is the least affected by systematic effects. To the statistical error a systematic one has to be added. It comes from the uncertainty in the overall mass scale due to the absolute calibration of the electromagnetic calorimeter, the uncertainty in the subtraction of uncorrelated low-momentum hadrons hitting the same calorimeter element, and the small inhomogeneities in the relative calibration of various cells of the calorimeter. It amounts to $\pm 3\%$.

The longitudinal momentum of the electron-neutrino system can be determined with a twofold ambiguity for the unmeasured component of the neutrino momentum. In fact, one does not know whether the electron was emitted forward or backward in the W rest frame. This results in two solutions for the fractional beam energy carried by the W, $x_W = p_L^W/p^{\max}$, and for the fractional longitudinal momenta of the quarks involved in the reaction, x_q and $x_{\bar{q}}$. The following relations hold (see Section 3.3)

$$x_q x_{\bar{q}} = m_W^2/s \quad (7.3)$$

$$x_q - x_{\bar{q}} = x_W.$$

One can compute [114, 115] the longitudinal momentum of the neutrino p_L^{ν} such that the mass of the electron-neutrino system is m_W :

$$m_W^2 = m_e^2 + 2[(p_L^{\nu 2} + p_T^{\nu 2})(p_L^{e 2} + p_T^{e 2} + m_e^2)]^{1/2} - 2p_L^e p_L^{\nu} - 2\vec{p}_T^e \vec{p}_T^{\nu}. \quad (7.4)$$

This equation is of second order in p_L^e . There are, therefore, two solutions which correspond to two quadrivectors for the W , and therefore two values of θ^* , the angle between the electron (positron) and the proton (antiproton) in the W rest frame. If the transverse momentum of the W is small compared to its longitudinal momentum, one finds

$$\sin \theta^* = p_T^e / m_W, \quad (7.5)$$

corresponding to two solutions of $\cos \theta^*$:

$$\cos \theta^* = \pm (1 - \sin^2 \theta^*)^{\frac{1}{2}}. \quad (7.6)$$

To resolve the ambiguity one uses conservation of energy and momentum (the unphysical solution has $x_W > 1$) and the measurement of the forward energy in the two beam arms setting upper limits on x_q and $x_{\bar{q}}$, leading again to only one solution. In about 70% of the events the ambiguity can be resolved. For the remaining events the two solutions are quite close. They correspond to W 's produced almost at rest with an electron emitted around 90° in the W rest frame. In those cases the lowest x_W value is taken. A complete list of the W -candidates can be found in Ref. 89. The x_W distribution in Fig. 47 is in agreement with the theoretical expectations. Also the x_q and $x_{\bar{q}}$ distributions (Fig. 48a,b) agree with the expected x distributions for quarks and antiquarks, respectively, in the proton and antiproton. In addition, the charge of the W -particle enables us to extract the u - and d -quarks distribution in the proton and antiproton separately [remember $(u\bar{d}) \rightarrow W^+$ and $(\bar{u}d) \rightarrow W^-$], assuming valence-quark dominance in the W -production. These distributions are shown in Fig. 49a,b.

As emphasized in Section 3.4, the weak coupling $(V - A)$ of the W to the $q\bar{q}$ and $e\nu$ system gives rise to a charge asymmetry. To test this prediction, one is forced to use only those events where the charge of the electron is unambiguously determined: 29 events (20 negative and 9 positive) remain. The correlation between the transverse energy of the charged lepton and its emission angle in the laboratory frame (Fig. 50) should give an indication of the asymmetry. In fact, one expects that the events populate preferentially the left side of the distribution, which indeed is observed. However, it has to be pointed out that this distribution may be affected by the transverse and longitudinal momenta of the W , or the u and d quarks. Therefore it is better to study the charge asymmetry in the W rest frame. For this one has to remove the events where the $\cos \theta^*$ is ambiguous. The angular distribution of the charged leptons should exhibit in this frame a $(1 + \cos \theta^*)^2$ behaviour. This is shown in Fig. 51, where the data are corrected for experimental acceptance and where the efficiency of solving the x_W ambiguity was taken into account. There is excellent agreement with the theoretical prediction, supporting strongly the $(V - A)$ coupling of the weak current.

Additional information may be extracted from the $\cos \theta^*$ distribution. In fact, as shown by Jacob [116], for a particle with arbitrary spin J one expects

$$\langle \cos \theta^* \rangle = \langle \lambda \rangle \langle \mu \rangle / J(J+1), \quad (7.7)$$

where $\langle \mu \rangle$ and $\langle \lambda \rangle$ are, respectively, the global helicity of the production system ($u\bar{d}$) and of the decay system ($e\bar{\nu}$). For $(V - A)$ coupling one has $\langle \lambda \rangle = \langle \mu \rangle = -1$, $J = 1$, leading to a value $\langle \cos \theta^* \rangle =$

0.5. For $J = 0$ one obviously expects $\langle \cos \theta^* \rangle = 0$. Experimentally, one finds $\langle \cos \theta^* \rangle = 0.4 \pm 0.1$, supporting both the $J = 1$ assignment and maximal helicity states at production and decay. We remember that the choice of the sign $\langle \mu \rangle = \langle \lambda \rangle = \pm 1$ cannot be made, i.e. $(V + A)$ and $(V - A)$ coupling both at production and decay cannot be resolved without an interference experiment. On the other hand, a helicity flip ($\langle \lambda \rangle = -\langle \mu \rangle$) is excluded.

Let us now turn to the total cross-section determination. With $p_T^e > 20$ GeV/c one selects 47 events. For the calculation of the cross-section the azimuthal cut is not necessary, since the electron energy does not need to be known very accurately. The $W \rightarrow \tau \nu_\tau$ contamination in this sample is estimated to be 2 ± 2 events. The overall acceptance is 0.65 obtained from:

- i) the $p_T^e > 20$ GeV/c cut (0.80),
- ii) the jet veto requirement within $\Delta\phi = \pm 30^\circ$ (0.96 ± 0.02),
- iii) the electron-track isolation requirement (0.90 ± 0.07),
- iv) the geometrical acceptance (0.94 ± 0.03).

This results in

$$(\sigma B)_w = 0.53 \pm 0.08 \text{ (stat.)} \pm 0.09 \text{ (syst.) nb,} \quad (7.8)$$

in excellent agreement with expectations [58].

7.1.2 The UA2 experiment

The sample of 225 events with at least one electron candidate having $p_T^e > 15$ GeV/c still contains, in addition to genuine electrons, fake electrons resulting from misidentified high- p_T hadrons or jets of hadrons [102, 103].

In the case of genuine electrons, because of lepton number conservation the event must contain either another electron of opposite charge (e.g. $Z^0 \rightarrow e^+e^-$) or a neutrino (e.g. $W \rightarrow e\nu$), which are also emitted, in general, with high p_T .

If the electron candidate is not genuine, but results from a misidentified high- p_T hadron or jet of hadrons, one expects another jet of high- p_T hadrons to be present in the same event at an approximately opposite azimuth. Such a configuration is characteristic of events containing high- p_T hadronic jets.

In order to examine the structure of these events, one uses the fine segmentation of the calorimeters to group into clusters all adjacent cells containing an energy deposition in excess of 400 MeV. Most clusters have low transverse momentum, as expected for the soft secondaries associated with the spectator constituents of the incident p and $\bar{p}$. To select preferentially high- p_T particles resulting from the same hard collision as that which produced the electron candidate, one sets a 3 GeV/c threshold on the cluster transverse momentum. This value removes most of the spectator particles; the probability to observe at least one cluster with $p_T \geq 3$ GeV/c in minimum-bias events is measured to be only 15%.

One finds that 45 events contain no additional high- p_T cluster, the electron candidate being the only high- p_T particle observed in the detector. Such events either contain a neutrino, as in the case of $W \rightarrow e\nu$ decays, or contain other high- p_T particles which have escaped detection because they were emitted outside the detector acceptance. Two-jet events, with one of the jets misidentified as an electron and the other undetected, are expected to contribute to the latter category. In this case, the p_T distribution of these fake electrons is expected to show the rapidly falling behaviour typical of the jet p_T spectrum.

The remaining 180 events contain at least one additional high- p_T cluster. In order to reduce the two-jet background in this sample, one considers for each event all the jets whose momentum $\vec{p}^{\text{jet}}$ is separated in azimuth from $\vec{p}^e$ by an angle $\Delta\phi > 120^\circ$; these jets are used to define the quantity

$$\rho_{\text{opp}} = -(\vec{p}_T^e \cdot \vec{p}_T^{\text{jet}}) / |\vec{p}_T^e|, \quad (7.9)$$

where $\vec{p}_T^e$ and $\vec{p}_T^{\text{jet}}$ are the transverse momenta of the electron and the jet.

All events which have $\rho_{\text{opp}} > 0.2$ are rejected (a total of 156 events including 8 Z^0 candidates). From the remaining 24 events, 16 events have $\rho_{\text{opp}} = 0$, i.e. no jet activity at all in an azimuthal sector $\pm 60^\circ$ wide opposite to the electron direction.

The sample of the 148 events is considered as a pure background sample. From these events it is now possible to estimate the background in the 45 events with no jet and in the 24 events with no opposite jet. This background is mainly coming from hadrons misidentified as electrons. As a result, the expected number of background events with $p_T^e > 25$ GeV/c is 0.6 events for the first sample and 0.9 events for the second sample.

The two event samples (45 and 24 events) are expected to contain electrons from $W \rightarrow e\nu$ decay. To study these events further one adds together their p_T^e distributions and the expected background contributions estimated as discussed above.

The p_T^e distribution of these 69 electron candidates is shown in Fig. 52 together with the background and signal estimate. This distribution shows an accumulation of events near $p_T^e \approx 40$ GeV/c, which is distinctive of the Jacobian peak expected for $W \rightarrow e\nu$ decay. For $p_T^e > 25$ GeV/c the expected background amounts to 1.5 events.

As a further check that the 37 electron candidates with $p_T^e > 25$ GeV/c are indeed consistent with being genuine, one applies the strict electron identification criteria (see Table 6) to the sample of Fig. 52. The p_T^e distribution of the surviving events is shown in Fig. 53. Above $p_T^e = 25$ GeV/c the number of events is now 33, which is consistent with that expected from the ratio of the cut efficiencies.

The presence of a neutrino in $W \rightarrow e\nu$ events can be detected by the resulting transverse momentum imbalance. For each event one calculates

$$\vec{p}_T^\nu = \vec{p}_T^W - \vec{p}_T^e, \quad (7.10)$$

where all vectors are projected onto a plane normal to the beams. The W transverse momentum $\vec{p}_T^W$ can be determined from the momenta of all other particles, or jets of particles, observed in the event in addition to the electron candidate:

$$\vec{p}_T^W = -(\sum \vec{p}_T^{\text{jet}} + \vec{p}_T^{\text{sp}}), \quad (7.11)$$

where the sum extends over all observed jets and $\vec{p}_T^{\text{sp}}$ is the total transverse momentum carried by the system of all spectator particles. In an ideal detector, $\vec{p}_T^{\text{sp}}$ is just the sum $\sum \vec{p}_T^{\text{sp}}$ over all spectator particles. However, corrections must be applied to take into account the errors in the measurement of the spectator transverse energies and the incomplete angular coverage of the apparatus.

Figure 54 shows the corrected distribution of p_T^W for the 37 $W \rightarrow e\nu$ candidates having $p_T^e > 25$ GeV/c. The mean value of p_T^W is $\langle p_T^W \rangle = 6.9 \pm 1.0$ GeV/c. There are

5 events with $p_T^W > 12$ GeV/c, the highest p_T^W value being 29.6 GeV/c. QCD predictions [67] are consistent with the observed distribution.

The neutrino transverse momentum $\vec{p}_T^\nu$ is obtained from Eq. (7.10) and one determines then its two orthogonal components p_T^ν and p_η^ν with respect to the electron direction. Figure 55a and 55b show all 69 events of the sample in the (p_T^e, p_T^ν) and (p_T^e, p_η^ν) planes, respectively. Most of the events populate the regions $p_T^\nu \approx p_T^e$, $p_\eta^\nu \approx 0$, where $W \rightarrow e\nu$ events with small values of p_T^W ($p_T^W \ll m_W$) are expected to fall. There are three events with $p_\eta^\nu > 20$ GeV/c, which are clearly separated from the rest of the population. The only event with $p_T^e > 25$ GeV/c is that having the highest p_T^W in the distribution of Fig. 54. This event is interpreted as the production of a high- p_T W recoiling against a jet, followed by the decay $W \rightarrow e\nu$.

With a maximum-likelihood method the mass of the W -particle, m_W , is determined from the p_T^e distribution in Fig. 53, using only the 37 events with $p_T^e > 25$ GeV/c. Assumption is made about the proton (antiproton) structure function [117] and the shape of the W transverse momentum [111]. Fixing the width of the W to $\Gamma_W = 2.7$ GeV, and including the contribution of (1.5 ± 0.1) background events and of (2.5 ± 0.9) misidentified Z^0 events (one electron not detected), one obtains the fitted value:

$$m_W = 83.1 \pm 1.9 \pm 1.3 \text{ GeV}/c^2. \quad (7.12)$$

where the systematic error is due to energy calibration uncertainties of the electromagnetic calorimeters.

Another method of mass determination consists in a fit done to the transverse mass m_T^W , by taking only those W -events with $p_T^W < 5$ GeV/c. The result is

$$m_W = 83.5 \pm 1.6 \pm 1.3 \text{ GeV}/c^2. \quad (7.13)$$

Let us now turn to the charge asymmetry. For this analysis, 11 events with $p_T^e > 15$ GeV/c in the forward regions can be used, because of the knowledge of the charge of the lepton. A comparison between the electron momentum p and the energy E , as measured in the calorimeter, is made in Fig. 56, which shows the events in the plane $(1/p, 1/E)$, where $p = qp \cos \theta / |\cos \theta|$, θ being the laboratory angle of the charged lepton and q its charge. A clear asymmetry is visible in Fig. 56, where most of the events lie on one side of the plot. In addition, one can define an asymmetry parameter

$$\alpha = \{[N(e^+) - N(e^-)] + [N(e^-) - N(e^+)]\} / N_{\text{tot}}, \quad (7.14)$$

where $N(e^\pm)$ is the number of $e^\pm$ for $142.5^\circ < \theta < 160^\circ$, $N(e^\pm)$ is the number of $e^\pm$ for $20^\circ < \theta < 37.5^\circ$, and N_{tot} is the total number of events. The observed asymmetry is

$$\alpha_{\text{exp}} = 0.82 \pm 0.17 \quad (7.15)$$

in agreement with the $(V - A)$ prediction of $\alpha = 0.67$ for $p_T^e > 15$ GeV/c.

For the 40 events the total cross-section is determined to be

$$(\sigma B)_W = 0.53 \pm 0.1 \pm 0.1 \text{ nb} \quad (7.16)$$

after subtraction of the Z^0 contribution and the $W \rightarrow \tau\nu_\tau$ contribution, and taking into account the detector acceptance.

7.2 The decay $W \rightarrow \mu\nu$ in the UA1 experiment

The scanning of the 689 events [101] rejected those muon candidates originating from cosmic rays, beam-gas interactions and leakage of hadronic showers or obvious interactions, showing multiple tracks in the muon chambers or large pulse height in the calorimeters. Parallel to it a selection program was applied, giving the same result as the scanning procedure. It was very helpful to reject the difficult background arising from $K \rightarrow \mu\nu$ decay. In this process the transverse kick to the muon from the decay balances the deflection in the magnetic field and therefore straightens the muon track. The result is a fake high-momentum muon. Asking for a good track fit removes a large part of this background. There remained 144 events. Next the events showing a jet activity had to be removed. They were of two kinds:

- i) Events with the muon in a jet, i.e. if the sum of the transverse momenta of other CD tracks around the muon in a cone of half width $\Delta r = (\Delta\phi^2 + \Delta y^2)^{\frac{1}{2}} < 0.7$ exceeded 3 GeV/c or if the UA1 jet algorithm [108, 109] found a calorimeter jet with $E_T > 10$ GeV/c for which the jet axis was within the above cone.
- ii) Events with a clearly visible jet within an azimuthal angle cone $|\Delta\phi| < 30^\circ$ opposite to the muon, as in the electron case (see Section 7.1.1).

A total of 107 events had these characteristics; they were removed and 37 events remained. Careful scanning removed still some cosmons and $K \rightarrow \mu\nu$ decays and reduced the sample to 18 events. A clean W-sample was obtained by applying a cut $E_T^\mu > 15$ GeV, leading finally to 14 $W \rightarrow \mu\nu$ candidates. A typical event is shown in Fig. 57.

Additional validation of the muon is done by studying the energy deposited in the calorimeters and the agreement of the extrapolated CD track with the track in the muon chamber. The ionization loss in the calorimeters was compared with results from ionization measurements with 45 GeV muons in a test beam (see Section 5.1.3). They are in good agreement. Also the muon momenta measured by the CD alone and with the muon chamber in addition (see Section 6.1.2) agree well (Fig. 58).

The most dangerous background, as discussed above, comes from the decay of medium-energy kaons leading to fake high-momentum muons. In addition, it fakes also a high-momentum neutrino, since momentum balance has to be preserved. With the help of a Monte Carlo calculation, using the experimentally known inclusive particle distribution ($3 < p_T < 15$ GeV/c) and simulating the detector, this background amounts to 0.9 events for $p_T^\mu > 15$ GeV/c. In addition, 5 events are expected from high-momentum ($p_T > 15$ GeV/c) pion or kaon decays. They are rejected by the selection criteria, since the apparent jet in those events balances transversely the high- p_T muon.

As in the electron case, the 14 W-candidates show a strong correlation between the transverse momentum of the neutrino and the transverse momentum of the muon, both in magnitude and direction (Fig. 59). This suggests a common origin for the muon and the neutrino, the W-particle. The muon transverse energy distribution (Fig. 60) shows an expected Jacobian peak, which is smeared out due to the large errors in momentum determination. However, the transverse momentum of the W, $\vec{p}_T^W = \vec{p}_T^\mu + \vec{p}_T^\nu$, is only affected by the errors in the calorimetry. Indeed, it is

$$\vec{p}_T^{\nu} = -\vec{p}_T^{\mu} + \vec{E}_T^{\text{cal}}, \quad (7.17)$$

where $\vec{E}_T^{\text{cal}}$ is the vector of the measured transverse energy in the event (we neglect the small energy deposited by the muon). Therefore

$$\vec{p}_T^W = \vec{E}_T^{\text{cal}}. \quad (7.18)$$

The distribution of $\vec{p}_T^W$ is shown in Fig. 61 in perfect agreement with the measurement from the $W \rightarrow e\nu$ sample, shown in Fig. 45. The two events with the highest $\vec{p}_T^W$ have a jet which locally balances the transverse momentum of the W.

The mass of the W-particle is again obtained by using a maximum-likelihood method for the transverse mass distribution, leading to a value of

$$m_W = 81^{+6}_{-7} \text{ GeV}/c^2, \quad (7.19)$$

in excellent agreement with the m_W value from $W \rightarrow e\nu$ decay (Sections 7.1.1 and 7.1.2). In the calculation of the cross-section an overall acceptance of 0.2 ± 0.035 comes into play. This value is composed of the geometrical efficiency (0.49), the efficiency of the quality cuts for the CD tracks (0.46 ± 0.07), and the correction for the jet veto and track isolation requirements (0.87 ± 0.07). The resultant cross-section is

$$(\sigma B)_W = 0.67 \pm 0.17 (\text{stat.}) \pm 0.15 (\text{syst.}) \text{ nb}, \quad (7.20)$$

where the systematic error includes the error on the acceptance and the 15% uncertainty in the integrated luminosity. Comparing it with the electron case (7.6) we obtain a ratio

$$R = (\sigma B)_\mu / (\sigma B)_e = 1.24^{+0.6}_{-0.4}, \quad (7.21)$$

in agreement with 1.

8. SEARCH FOR Z^0

In the same data sample as for the W case, the UA1 and UA2 groups have searched for the neutral boson Z^0 .

8.1 The decay $Z^0 \rightarrow e^+e^-$

8.1.1 The UA1 experiment

Again the data from the express line have been taken to perform this search [45, 118, 119]. This time only events with two contiguous cells of the electromagnetic calorimeters with $E_T > 25$ GeV and a small (or no) energy deposition (≤ 600 MeV) in the hadron calorimeter immediately behind them were retained. Only 225 events from the initial 1.5×10^5 events satisfied these criteria (Fig. 62). Asking for a matching high-transverse-momentum track in the CD for at least one cluster reduces the sample to 129 events. This is further reduced to 13 events if the isolation of the track is required. If the last cut is also applied to the second cluster, only 4 (e^+e^-) events remain, which have

a common value of invariant mass. The main parameters are given in Table 7. One of these events is shown in Fig. 63. A common feature of all events is the two large electromagnetic energy depositions, as illustrated in Fig. 64, and the absence of missing transverse energy, indicating no evidence for the emission of energetic neutrinos.

The electrons have been validated according to their energy deposition in the gondolas, in agreement with the distribution obtained from the $W \rightarrow e\nu$ events and from the results of the test-beam measurements. Furthermore, the correctness of the electron assignment can be obtained by comparing the impact of the electron tracks on the gondolas, as measured by the central detector, with the centroid of the energy deposition measured from the ratio of the photomultiplier signals (Fig. 65).

Comparing the momenta measured in the CD with the energy deposition in the gondolas (Fig. 66), one finds perfect agreement for 6 of the 8 tracks. One track is unmeasurable because it travels less than 15° parallel to the magnetic field; another one shows a clear inconsistency between energy and momentum determination (see Table 7). In fact, the momentum measured by curvature in the CD is (9 ± 1) GeV/c, much smaller than the corresponding calorimeter deposition of (48 ± 1) GeV. This event is interpreted as the likely emission of a “hard” photon accompanying the electron. A photon is defined as having a transverse energy deposition in excess of 10 GeV in two adjacent electromagnetic calorimeter cells, with no associated charged tracks with transverse momentum in excess of 1 GeV/c, and less than 600 MeV/c energy deposited in the hadronic calorimeter cells immediately behind the electromagnetic cluster. In addition, it is required that this cluster is not part of a large calorimeter jet structure as defined by the UA1 jet algorithm [108, 109]. Subsequent calibrations of the gondolas indicated that the centroid of the energy deposition in the calorimeters was significantly displaced with respect to the incident electron track, indicating an azimuthal angle of $\Delta\phi = (14 \pm 4)^\circ$ between the charged and neutral components [120]. This makes the possibility of an external bremsstrahlung in the vacuum pipe and detector window seem unlikely. The estimated probability of an external bremsstrahlung exceeding the angle and energy observed is [121] about 0.005, or 2% for the sample of four events.

Despite this peculiar effect—with which we will be faced again later—let us evaluate the possible backgrounds. Trivial sources of background are not expected to exhibit such a clustering at high masses. No $(e\mu)$ pair and no same-sign dilepton has been observed at such high masses, indicating that the background is probably small. Nevertheless, several possible sources of background have been considered [45].

- i) Two ordinary large-transverse-momentum jets, both fragmenting into isolated high-momentum tracks, may simulate electrons. Using the measured jet cross-section and jet charged-fragmentation function [109] and requiring the electromagnetic signature for the isolated tracks the background ($\sim 2 \times 10^{-3}$ events) becomes negligible.
- ii) Heavy-flavour jets may decay into leptons. This background was estimated from events which have a high-transverse-momentum lepton accompanied by a recoil jet to be $< 10^{-3}$ events.
- iii) Production of lepton pairs by the Drell–Yan process. With a cross-section of approximately 10^{-5} nb/GeV [54, 66, 67] for a dilepton mass of 80 GeV, this yields $\sim 2.5 \times 10^{-2}$ events for masses greater than 80 GeV. When requiring a threshold of only 15 GeV for each electromagnetic cluster, one (e^+e^-) pair appears with a mass just above 30 GeV/c², in agreement with the expectation from the Drell–Yan process.
- iv) W^+W^- pair creation is expected to be entirely negligible at our energy [122].

- v) Onium decay from a new heavy quark is estimated [123] to be much too small to account for the desired effect in the mass range of $95 \text{ GeV}/c^2$.

Therefore, one can safely conclude that there is no background producing the number and features of the observed events. All the observations are in agreement with the hypothesis that events are due to the production and decay of the neutral intermediate vector boson Z^0 .

Assuming a Breit–Wigner distribution, one obtains a mass

$$m_{Z^0} = 95.6 \pm 1.4 (\text{stat.}) \pm 2.9 (\text{syst.}) \text{ GeV}/c^2, \quad (8.1)$$

where the systematic error reflects the uncertainty on the absolute calibration. The width of the Z^0 -peak is determined to be

$$\Gamma_{Z^0} < 8.5 \text{ GeV} \quad (90\% \text{ CL}). \quad (8.2)$$

With the known acceptance one evaluates the cross-section to be

$$(\sigma B)_{Z^0 \rightarrow e^+e^-} = 0.05 \pm 0.02 (\text{stat.}) \pm 0.009 (\text{syst.}) \text{ nb} \quad (8.3)$$

in agreement with expectation.

8.1.2 The UA2 experiment

After having applied the strict cut 1 (see Table 6) to both charged leptons, 24 events remain (see Section 6.2.2), whose invariant mass distribution [46, 102] is shown in Fig. 67. There are 12 events with both clusters in the central region, 8 events with one cluster in the central region and the other in the forward regions, and 4 events with both clusters in the forward regions. The effect of the cuts 2 to 6 in Table 6 has been studied with the initial Z^0 -sample. It turns out that the rejection power of any of these cuts acts quadratically if applied to both clusters simultaneously and that the shape of the mass distribution shows no dependence, within statistics, on the combination of cuts used. These two observations provide a simple method to estimate the background contribution from two-jet events to any mass region in a sample of events surviving a given combination of cuts.

Eight events of Fig. 67a survive if at least one cluster satisfies all the electron identification criteria. They are shown in Fig. 67b and they have a common value of invariant mass clustered around $90 \text{ GeV}/c^2$. Their relevant parameters are given in Table 8. For clusters passing cut 3 the cluster energy has been corrected to account for the calorimeter response as a function of the electron incidence angle and impact point. With this new energy the invariant mass values in Fig. 67 have been calculated.

The background to the eight events has been estimated in two ways:

- i) By comparing the event populations above and below $80 \text{ GeV}/c^2$ one finds an upper limit (90% CL) of 0.32 background events to the signal in Fig. 67b. This allows for 2.3 events with $m < 80 \text{ GeV}/c^2$ in Fig. 67b. None is observed.
- ii) One uses a much larger event sample by releasing cut 1 on one of the two clusters and evaluates the rejection power of cuts 2 to 6 on the other cluster. Absence of correlations between the fragmentation of the two jets in background events legitimates this procedure. The background estimated in this way amounts to 0.03 events.

If one applies all the electron identification criteria to both clusters, only events A, B, and C (see Table 8) survive, with an estimated background of 2×10^{-4} events above a mass of $80 \text{ GeV}/c^2$. Let us consider them in detail.

Events A and B are interpreted as clear examples of $Z^0 \rightarrow e^+e^-$ decays. Figure 68 shows the event A and the deposited energies of the electrons. However, event C consists of two electrons and a well-separated high-energy photon, comparable to the UA1 event. The estimated probability that such an event occurs with the observed features is 0.005, compared to normal $Z^0 \rightarrow e^+e^-$ decays. The event is shown in Fig. 69.

Let us turn to events D to H, characterized by the failure of one of the two electron candidates on the application of at least one of the strict selection criteria:

- i) Event D: The forward electron candidate is associated with a track measured in the vertex detector as pointing to a coil of the magnet. This configuration is consistent with the hypothesis of an electron initiating an electromagnetic shower in the magnet coil and the mass has been calculated accordingly.
- ii) Event E: The central electron candidate fails only cuts 4 and 5. Its associated C5 cluster has a charge equivalent to only 2.4 minimum ionizing particles and is (9 ± 2) mm away from the track impact, but this is still compatible with an electron-pair hypothesis. On the other hand, the forward electron is accompanied by another particle with $p = 2.3 \text{ GeV}/c$ hitting the same calorimeter cell. From minimum-bias events the estimated probability for such a configuration is 0.2%.
- iii) Event F: One of the central electron candidates fails only cut 3. The measured ratio of the signals from the two light-guides of the impact cell is 0.71 ± 0.02 instead of 0.88 ± 0.02 as predicted from the track impact. Careful inspection of the behaviour of the calorimeter cell during calibration and comparison of the above ratio with expected values from minimum-bias events do not help to clarify the situation. Only advancing the hypothesis of an additional neutral particle (a hard bremsstrahlung photon) hitting the edge of the same calorimeter cell explains the event. However, the absence of a C5 cluster facing this region would imply that the photon did not convert in the 1.5 radiation lengths thick converter.
- iv) Event G: One of the central electron candidates again fails cut 3 because additional energy ($\sim 3 \text{ GeV}$) is deposited in the neighbouring cell. This is consistent with the present knowledge of the detailed response of the central electromagnetic calorimeter, concerning lateral and longitudinal leakages. No tracks point to the cell and no C5 cluster faces it.
- v) Event H: Here the same is true as for event G but in addition the cut 4 is failed. This is because the electron candidate falls in a small region where C5 is non-operational.

The UA2 group comes to the following conclusion:

- i) Eight events (A to H) are observed with at least one electron passing the strict identification criteria. Their masses cluster in the $90 \text{ GeV}/c^2$ region with an expected background of only 0.03 events.
- ii) Five of these events (A to E) are either identified as, or perfectly compatible with, e^+e^- or $e^+e^-\gamma$ configurations.
- iii) The three remaining events (F to H) have always one cluster failing cut 3. Although this is clearly interpreted as an association of an electron with the cluster, the above-mentioned inconsistency has to be clarified in additional tests in an electron beam before drawing any definite conclusion.

The difficulties with events F to H led to a reconsideration of the background below the mass peak in Fig. 67b, this time owing to a possible contamination coming from quark-quark decays of the IVBs. However it is difficult to retain such a hypothesis because the sample in Fig. 67a contains only two extra events, compatible with an expectation of 0.7 events for $m > 80 \text{ GeV}/c^2$.

The most likely interpretation of the eight events is that they all result from the decays $Z^0 \rightarrow e^+ e^-$ or $Z^0 \rightarrow e^+ e^- \gamma$. On the other hand, for the determination of the mass and the width, only events A, B, C, and E can be used with confidence. One finds

$$m_{Z^0} = 92.7 \pm 1.7 (\text{stat.}) \pm 1.4 (\text{syst.}) \text{ GeV}/c^2, \quad (8.4)$$

where the systematic error accounts for the uncertainty on the overall energy scale. The corresponding width is evaluated to be

$$\Gamma_{Z^0} < 6.5 \text{ GeV} \quad (90\% \text{ CL}). \quad (8.5)$$

8.2 The decay $Z^0 \rightarrow \mu^+ \mu^-$ in the UA1 experiment

The search for this process was originally made from the express-line events [45, 124]. By requiring the same criteria as for the muons for the $W \rightarrow \mu \nu$ search, 30 dimuon candidates were left. Careful scanning of these events resulted in a total of 11 ($\mu^+ \mu^-$), 2 ($\mu^- \mu^-$), and 1 ($\mu^+ \mu^+$) candidates, the rest being cosmic-ray muons or leaking particles from showers. From these events, 10 have a dimuon mass less than $25 \text{ GeV}/c^2$ (Fig. 70), whereas 4 show a mass $m > 80 \text{ GeV}/c^2$. By selecting 2 muons from the filter line, the 14 events discussed above reappeared, but in addition, another event with two high-transverse momentum muons and a mass $> 70 \text{ GeV}/c^2$ was found. The reason for its absence in the express line lies in the fact that one muon in the event falls into a region where the dimuon trigger was not active. The parameters of the five dimuon events with high mass are listed in Table 9. Event F is displayed in Fig. 71.

The muons in the five high-mass events fulfil perfectly the criteria for muon identification. In fact, the energy losses in the calorimeters traversed by the muons are well within expectations of ionization losses of high-energy muons (Fig. 72). Also the coordinates and angles measured at the exit of the iron absorber are in complete agreement with the extrapolated CD tracks, as illustrated in Fig. 73, where the expectation from cosmic-ray muons is also shown.

The momentum of the muons is determined by three methods. The CD measures the momentum by magnetic deflection of the track. An independent measurement is obtained from the deviation of the muon from a straight line extrapolating the original direction at the vertex to the muon chambers (see Section 6.1.2). Finally, and this is only possible in the Z^0 case, an overall fit of the event determines the muon momenta. Since the missing transverse energy is small (see Table 9), indicating the absence of neutrino emission, the recoil of the hadronic debris must balance the transverse momentum of the $\mu^+ \mu^-$ pair. The directions of the two muons then suffice to determine their momenta, which are now only dominated by the errors in calorimetry.

We next discuss the five events in detail:

- i) Events F, H, and I are interpreted as $Z^0 \rightarrow \mu^+ \mu^-$ decays. Whereas events H and I show a minimum-bias behaviour for the hadronic debris, event F has a significant jet which balances the $\mu^+ \mu^-$ pair in transverse momentum.

- ii) Event G has in addition to the high-transverse-momentum muons a well-separated high-energy e.m. cluster, similar to the two events found in the e^+e^- search. Now, however, the bremsstrahlung nature of a photon is highly improbable. The event is interpreted as a $Z^0 \rightarrow \mu^+\mu^-\gamma$ decay. For the determination of the muon momenta, one does not apply the E_T balance, since the origin of such an event is unknown.
- iii) Event E, although having two high-momentum muons, has a dimuon mass which is 1σ away from the determined Z^0 -mass in the (e^+e^-) case. Nevertheless, the interpretation that this event is a $Z^0 \rightarrow \mu^+\mu^-$ decay is retained, because it is difficult to believe in any background larger than 10^{-2} events for a mass exceeding $75 \text{ GeV}/c^2$, as described in Section 8.1. In addition, this event shows a large jet activity, which balances perfectly the Z^0 in transverse momentum.

We conclude that the most likely interpretation of these five events is that they result from $Z^0 \rightarrow \mu^+\mu^-$ or $Z^0 \rightarrow \mu^+\mu^-\gamma$ decays. The mass of the 12 $(\mu^+\mu^-)$ events is shown in Fig. 74 together with the (e^+e^-) mass from the UA1 and UA2 experiments. By excluding the $\mu^+\mu^-\gamma$ event, one obtains a mass value of

$$m_{Z^0} = 85.6 \pm 6.3 \text{ GeV}/c^2, \quad (8.6)$$

which, combined with the value obtained for the $Z^0 \rightarrow e^+e^-$ events [see Eq. (8.1)], leads to an average value for the Z^0 -mass of

$$m_{Z^0} = 93.5 \pm 2.9 \text{ GeV}/c^2. \quad (8.7)$$

The transverse-momentum distribution of all Z^0 events, shown in Fig. 75, is in agreement with the QCD calculations. For the cross-section one obtains with the known acceptance

$$(\sigma B)_{Z^0 \rightarrow \mu^+\mu^-} = 0.105 \pm 0.050 \pm 0.015 \text{ nb}. \quad (8.8)$$

This, together with (8.3), determines an overall cross-section for the decay $Z^0 \rightarrow \ell^+\ell^-$ of

$$(\sigma B)_{Z^0 \rightarrow \ell^+\ell^-} = 0.071 \pm 0.024 \pm 0.013 \text{ nb}, \quad (8.9)$$

in agreement with the prediction of 0.055 nb [64], where a k factor of 1 is assumed.

Finally let us briefly comment on the low-mass dimuon events. The unlike-sign events $(\mu^+\mu^-)$ might be Drell–Yan pairs or stem from the semileptonic decay of quark–antiquark pairs which have been produced in QCD processes or come from a Z^0 decay. The origin of the like-sign events $(\mu^-\mu^-)$ or $(\mu^+\mu^+)$ might be *either* the hadronic decay of the W with subsequent semileptonic decay of the quarks, *or* the decay of a quark–antiquark pair, such as

$$\begin{array}{c} b\bar{b} \\ \downarrow \text{ } \mu^+\nu\bar{c} \\ \text{ } \rightarrow cX \\ \downarrow \text{ } \mu^+\nu s \end{array} \quad (8.10)$$

Another interesting possibility is the production of a $b\bar{b}$ pair with subsequent $B_0-\bar{B}_0$ mixing [125]. The background due to π or K decay has been evaluated, using the observed inclusive muon sample, to be $\leq 10^{-2}$. A detailed study of the observed events can be found elsewhere [124].

9. CONCLUSIONS AND OUTLOOK

9.1 Comparison with the $SU(2) \times U(1)$ model

Combining the information on the $W^\pm$ and Z^0 obtained by the two experiments allows a direct comparison with the $SU(2) \times U(1)$ standard model [23, 24].

In Table 10 are summarized all the experimental results related to the IVBs, $W^\pm$ and Z^0 [45, 46, 99, 101, 103, 124]. For the comparison with the model we will establish an average between the UA1 and UA2 mass determinations, taking the systematic errors into account [126]. We obtain

$$m_W = 82.1 \pm 1.7 \text{ GeV}/c^2 \quad (9.1)$$

$$m_{Z^0} = 93.0 \pm 1.7 \text{ GeV}/c^2 \quad (9.2)$$

$$\Delta m = m_{Z^0} - m_W = 10.9 \pm 1.6 \text{ GeV}/c^2. \quad (9.3)$$

The error in the mass difference is smaller than that in the separate masses because of energy scale uncertainties which cancel in the difference. Figure 76 represents the experimental and theoretical situation [127]. One can now extract the renormalized value of $\sin^2 \theta_W$ at mass scale m_W following Eq. (3.3):

$$\sin^2 \theta_W^R = (38.5/m_W)^2. \quad (9.4)$$

Inserting the value of m_W from Eq. (9.1) one finds

$$\sin^2 \theta_W^R = 0.220 \pm 0.009, \quad (9.5)$$

in excellent agreement with the renormalized value of $\sin^2 \theta_W^R = 0.215 \pm 0.014$ [Eq. (3.4)], deduced from neutral-current experiments. Using the information on the Z^0 -mass, one can determine the parameter ρ , related immediately to the isospin of the Higgs particle:

$$\rho = m_W^2/m_{Z^0}^2 \cos^2 \theta_W^R. \quad (9.6)$$

Using the values (9.1), (9.2), and (9.5), one finds

$$\rho = 1.000 \pm 0.036, \quad (9.7)$$

in perfect agreement with the prediction of $\rho = 1$ for a Higgs doublet. Let us point out that ρ deviates from 1 at most by 3%, owing to radiative corrections involving possible new fermion generations [128]. The present value (9.7) seems to indicate no such new fermion families.

We conclude that, within errors, the observed experimental values are completely compatible with the $SU(2) \times U(1)$ model [23, 24].

In addition we point out that very massive W -like objects in addition to the observed $W^\pm$, postulated for instance in right-left-symmetric models [129], are ruled out. In fact, no event with p_T^e or p_T^ν in excess of the expected distributions for $W \rightarrow e\nu$ events has been observed (see Figs. 41 and

42) in the UA1 experiment. This sets a limit on the cross-section at $(\sigma B)_W \leq 30$ pb at 90% CL, corresponding to a mass $m_W > 170$ GeV/c², if standard couplings and quark distributions are used to evaluate the cross-sections.

9.2 Number of neutrinos

As already discussed in Section 3.2, each neutrino contributes with 180 MeV to the total decay width of the Z^0 , assuming $m_\nu \ll m_Z/2$.

From the measured decay width Γ_{Z^0} , the UA1 experiment sets a limit of $N_\nu < 31$ [119] and the UA2 experiment $N_\nu < 22$ [103]. An additional determination of N_ν can be made by comparing the leptonic branching ratios of $W^\pm$ and Z^0 . The following relation holds:

$$(\sigma B)_{Z \rightarrow e^+e^-} / (\sigma B)_{W \rightarrow e\nu} = (\sigma_Z / \sigma_W) (\Gamma_{Z \rightarrow e^+e^-} / \Gamma_{W \rightarrow e\nu}) (\Gamma_W^t / \Gamma_Z^t), \quad (9.8)$$

where B is the branching ratio, Γ the partial width, and Γ^t the total width. Equation (9.8) has the advantage that the ratio of the cross-sections σ_Z / σ_W appears in it. This value is known to be 0.24 ± 0.05 [58] and is much more reliable than the individual cross-section values since many of the uncertainties of the QCD models may cancel. With the help of a Monte Carlo calculation the UA1 group has determined $(\sigma B)_Z / (\sigma B)_W$ and Γ_Z^t , and the maximum likelihood fit gives [119]

$$(\sigma B)_Z / (\sigma B)_W = 0.098^{+0.052}_{-0.035} \quad (9.9)$$

and

$$\Gamma_Z^t = 1.8^{+5.5}_{-1.1} \text{ GeV}. \quad (9.10)$$

This is illustrated in Fig. 77 together with one-sided 90% and 95% confidence limits. Also shown are the values expected for different numbers of neutrino families, where the hashed areas just reflect the 20% error in σ_Z / σ_W . From this plot one can conclude that

$$N_\nu < 18 \quad (90\% \text{ CL}). \quad (9.11)$$

If one fits simultaneously the ratio and the mass distribution, an upper limit of $N_\nu < 14$ is obtained at 90% CL [130]. This limit is by far the best one today. Measurements give limits of $N_\nu \lesssim 1000$ for $K \rightarrow \nu\nu$ [131] and $N_\nu \lesssim 100$ for $e^+e^- \rightarrow \gamma\nu\nu$ [132]. A limit $N_\nu < 4$ (or $N_\nu > 6000$) is obtained from cosmological nucleosynthesis calculations [133].

9.3 Z^0 anomalies

The three events interpreted as $Z^0 \rightarrow \ell^+ \ell^- \gamma$ decay may be a possible indication for new physics. A summary of the main parameters of these events is given in Table 11. The dominant feature is the high photon energy and the large opening angle between the charged lepton and the photon. As already mentioned the possibility of a bremsstrahlung with such an energy and angle observed is very small.

Different theoretical explanations, such as composite weak bosons [134] or excited leptons [135], have been postulated. Therefore the UA1 group searched for analogous phenomena in the

charged IVB channels, namely for processes leading to final states of the type $e\nu\gamma$ and $\mu\nu\gamma$, where γ is a hard photon with $E_T^\gamma > 10 \text{ GeV}$ [120]. No evidence was found for such events. This places an upper limit on R , defined as

$$R = \text{BR}(X \rightarrow \ell\nu\gamma) / \text{BR}(W \rightarrow \ell\nu), \quad (9.12)$$

provided one adopts a model for the decay kinematics of the process. The limits for R were evaluated for three cases:

- i) Three-body phase space: $X = W \rightarrow \ell\nu\gamma$. The value $R \leq (90\% \text{ CL})$ is found.
- ii) Excited lepton:

$$\begin{array}{c} X = W \rightarrow \ell^* \nu \\ \quad \quad \quad \downarrow \\ \quad \quad \quad \ell\gamma \end{array}$$

The upper limit on R is 0.2 at a mass of $m_{\ell^*} = 10 \text{ GeV}/c^2$, falling to 0.1 at $70 \text{ GeV}/c^2$ (95% CL).

- iii) Radiative decay of a heavy W into a lighter W :

$$\begin{array}{ccc} X = \bar{W} \rightarrow W\gamma & \text{or} & X = W \rightarrow \bar{W}\gamma \\ \quad \quad \quad \downarrow & & \quad \quad \quad \downarrow \\ \quad \quad \quad \ell\nu & & \quad \quad \quad \ell\nu \end{array}$$

Here R varies also from 0.2 to 0.1 (90% CL) going from low to high $\bar{W}$ -masses. The result for R from the last two processes is shown in Fig. 78.

A similar search for $e\nu\gamma$ events has been undertaken by the UA2 group [103]. There is no evidence for an anomalous $W \rightarrow e\nu\gamma$ signature.

It is certainly premature to draw any conclusion about the origin of the $\ell\nu\gamma$ events; only future $p\bar{p}$ runs may clear up this situation.

9.4 The future of $p\bar{p}$ physics

The possible new phenomena encountered in the last section are certainly one goal to follow up in the next runs. Therefore, let us end by discussing the physics that can be done in the future.

Firstly, the search for the sixth quark, the top quark, is the most interesting challenge. In fact, the decay of the W into (tb) pairs constitutes a significant source of top quarks, provided the mass of the top quark is lower than the kinematical limit due to the W -mass. Owing to the colour factor the yield of $W \rightarrow tb$ decays is three times larger than in the leptonic decays of the W . The top quark itself must be identified through its decay. Here the semileptonic decay ($\sim 10\%$) into electrons or muons seems to give the best signature. However, the transverse momenta of the leptons are on the average smaller than $10 \text{ GeV}/c$ for top-quark masses ranging from 30 to $40 \text{ GeV}/c^2$, so that the small momentum tail is affected by background. With the present luminosity few events are expected [65] and the analysis for the above decay is still under way. Additional information about the decay vertices, extension of the muon trigger to small θ angles, and higher luminosity will certainly clear the situation on the top quark. Let us add here that also diffractively produced top quarks and $t\bar{t}$ production through gluon-gluon interaction might reveal the top signal, provided the top mass is not too high and the cross-section therefore not too low.

Secondly the search for the Higgs particle is most important. The properties of this particle are only elusively defined. The rate for its production and the experimental method to search for it

are strongly dependent on the mass, which is unrelated to any observable parameter. There are two very clear ways of looking for the Higgs particle:

- i) the decay of an IVB into a Higgs particle and a fermion–antifermion pair, such as $W^\pm \rightarrow H^0 + f + \bar{f}$;
- ii) the associate production of an IVB and of a Higgs particle mediated by a virtual boson propagator, such as $p\bar{p} \rightarrow W^\pm + H^0 + X$.

Both methods require, in addition to the presence of a virtual or real IVB, a positive identification of the decays of the H^0 , which is expected to decay preferentially into a pair of the most massive fermions kinematically accessible. The invariant mass distribution of the fermion–antifermion pair and the transverse momenta distribution of the IVBs are significantly different in both cases and can be used as an additional signature of the two processes. The branching ratio for the production of the Higgs particle, compared to normal IVB production, ranges from 10^{-3} to 10^{-4} for a mass of the Higgs particle of 10 GeV/c² and 40 GeV/c².

With improved resolution of electromagnetic calorimetry and high statistics, the $W^\pm$ and Z^0 parameters can be determined more precisely in order to test the standard model. The prediction on the QCD corrections to the masses are precise, so that very small deviations would yield an unambiguous signal for unanticipated effects. In other words, the deviation from $\rho = 1$ may indicate new heavy fermions or additional scalar weak multiplets in the electroweak theory. In addition, the charge asymmetry of the $W^\pm$ and Z^0 can be accurately determined and provides important information on the V,A couplings of charged and neutral currents. It may also indicate the presence of new, sequential heavy leptons. Finally, the number of lepton families can be determined by measuring the natural width of the Z^0 . Remember each neutrino increases this width by 180 MeV (see Section 3.2). The new particles which theorists predict, such as the supersymmetric [136], technicolour [137], and new W and Z particles, will affect the widths of the IVBs, provided their mass is substantially lighter than 80 GeV/c².

In order to cope with all this interesting physics, an upgrading program has been proposed by the UA1 group [138], which consists in replacing the existing electromagnetic calorimetry by a uranium calorimeter filled with tetramethylsilane, in adding a microvertex detector, in extending the acceptance and trigger capability of the forward muon detector by adding more shielding and chambers, and in instrumenting the iron shield with Iarocci tubes. The UA2 group plans to replace their forward detectors by new calorimeters in order to reach a better homogeneity, which helps in the energy determination of the events. Parallel to it a new antiproton accumulator, the ACOL [139], will be built, so that luminosities of $> 10^{30} \text{ cm}^{-2} \text{ s}^{-1}$ can be reached.

With all these innovations it is hoped to collect data for an integrated luminosity of 10^4 nb^{-1} in the coming years. This will provide a large number of IVBs, e.g. about 6000 $W \rightarrow e\nu$ events, and should contribute a lot to the understanding of elementary particles before the new machines, Tevatron, SLC, LEP and HERA, become operational.

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combined if their separation in phase space satisfies the cut $\Delta R = [(\Delta\eta)^2 + (\Delta\phi)^2]^{1/2} < 1$ ($\Delta\phi$ in radians). The remaining soft particles are added to the nearest jet in $\Delta\eta$ and $\Delta\phi$, provided the relative p_T is < 1 GeV and $\Delta\theta < 45^\circ$. A jet is considered valid if $E_T^{\text{jet}} > 10$ GeV. This same procedure is used for central-detector tracks with appropriately adjusted parameters (see also Ref. 109).

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Table 1
AA design parameters

Central momentum	3.5 GeV/c
Circumference	151.08 m at injection (CPS/4) 155.84 m at stack centre
Working point	$Q_x = 2.284, Q_y = 2.276$, at stack centre
Total $\Delta p/p$	$\pm 30 \times 10^{-3}$
Vacuum pressure	10^{-10} Torr
Target	Tungsten rod 110 mm long, 3 mm $\varnothing$, followed by magnetic horn
Protons per pulse on target	1×10^{13} every 2.4 s
Antiprotons per pulse	2.5×10^7 in AA acceptance
Acceptance (injection)	100 hor., 100 vert. $\pi \cdot \text{mm} \cdot \text{mrad}$ at $\Delta p/p = 7.5 \times 10^{-3}$
(transfer to CPS)	7.6 hor., 4.5 vert. $\pi \cdot \text{mm} \cdot \text{mrad}$ at $\Delta p/p = 1.1 \times 10^{-3}$
First fill ($1 \times 10^{12} \bar{p}$)	50,000 pulses, 33 h
Refill ($6 \times 10^{11} \bar{p}$)	30,000 pulses, 20 h

Table 2
SPS performance as a collider in 1982 and 1983

	1982	1983
AA stacking rate ($10^9 \bar{p} \text{ h}^{-1}$)	5.5	6.6
Transfer efficiency (%)	70	75
Proton intensity per bunch (10^{11})	1	1.4
Antiproton intensity per bunch (10^{11})	0.13	0.15
Normalized emittances ($10^{-6} \pi \cdot \text{rad} \cdot \text{m}$): protons	20	18
antiprotons	15	14
Beam-beam parameter ξ	0.0025	0.004
$\beta_H^* (\text{m}), \beta_V^* (\text{m})$	2×1	1.3×0.65
Luminosity ($10^{29} \text{ cm}^{-2} \text{ s}^{-1}$)	0.53	1.6
Lifetime (h)	18	16
Integrated luminosity per day ($\text{nb}^{-1} = 10^{-33} \text{ cm}^{-2}$)	2.2	6.2

Table 3
General characteristics of the central detector

Gas mixture	Argon (40%) + ethane (60%)
Drift field and gap length	1.5 kV/cm, 18 cm
Drift velocity	5.3 cm/ μ s
Drift angle	23° at $ \vec{B} = 0.7$ T
Anode plane arrangement:	
a) Distance between sense wires	10 mm
b) Wire length	80 cm min., 220 cm max.
c) Sense wire characteristics	35 μ m Ni-Cr stretched at 80 g
d) Field wire characteristics	100 μ m gold-plated Cu-Be stretched at 200 g
Cathode plane structure:	
a) Distance between wires	5 mm
b) Wire characteristics	150 μ m gold-plated Cu-Be stretched at 200 g
Total number of wires	22,800
Total number of sense wires	6110

Table 4
Calorimetry

Calorimeter		Angular coverage θ (°)	Thickness		Cell size		Sampling step	Segmentation in depth	Resolution
			No. rad. lengths	No. abs. lengths	$\Delta\theta$ (°)	$\Delta\phi$ (°)			
Barrel	e.m.: gondolas	25–155	26/ $\sin\theta$	1.1/ $\sin\theta$	5	180	1.2 mm Pb 1.5 mm scint.	3.3/6.6/9.9/6.6 X_0	0.15/ $\sqrt{E}$
	hadr.: C's		–	5.0/ $\sin\theta$	15	18	50 mm Fe 10 mm scint.	2.5/2.5 λ	0.8/ $\sqrt{E}$
End-caps	e.m.: bouchons	5–25	27/ $\cos\theta$	1.1/ $\cos\theta$	20	11	4 mm Pb 6 mm scint.	4/7/9/7 X_0	0.12/ $\sqrt{E_T}$
	hadr.: I's	155–175	–	7.1/ $\cos\theta$	5	10	50 mm Fe 10 mm scint.	3.5/3.5 λ	0.8/ $\sqrt{E}$
Calcom	e.m.	0.7–5	30	1.2	4	45	3 mm Pb 3 mm scint.	4 $\times$ 7.5 X_0	0.15/ $\sqrt{E}$
	hadr.	175–179.3	–	10.2	–	–	40 mm Fe 8 mm scint.	6 $\times$ 1.7 λ	0.8/ $\sqrt{E}$
Very forward	e.m.	0.2–0.7	24.5	1.0	0.5	90	3 mm Pb 6 mm scint.	5.7/5.3/5.8/7.7 X_0	0.15/ $\sqrt{E}$
	hadr.	179.3–179.8	–	5.7	0.5	90	40 mm Fe 10 mm scint.	5 $\times$ 1.25 λ	0.8/ $\sqrt{E}$

Table 5

General characteristics of the muon-chamber detector

Gas mixture	Argon (40%) + ethane (60%)
Drift field	$> 0.5 \text{ kV cm}^{-1}$
Drift space	7 cm
Drift velocity	5 cm/ μs
Sense-wire length	555 cm max.
Sense-wire characteristics	50 μm stainless-steel stretched at 300 g
Cathode characteristics	0.2 mm thick copper strips, 1.8 cm wide for -5 kV and 2.2 cm wide for -7 kV
Total number of sense wires	5180

Table 6
Electron identification criteria (loose cuts)*)

Requirements	Central region		Forward regions	
	Description	$\eta(\%)$	Description	$\eta(\%)$
1. Presence of a calorimeter cluster (preselection)	Electromagnetic cluster size $R_{\theta,\phi} < 0.5$ Hadronic leakage $H < 0.035 + 0.051 \ln E_{cl}$	99	Cluster size ≤ 2 cells Energy fraction in adjacent cells $< 5\%$ Hadronic leakage $H < 0.02$	99
2. Presence of a track	Reconstructed in the vertex detector in both transverse and longitudinal projections	90	Reconstructed a) in the vertex detector in transverse projection with $N \geq 1$ signal in the two inner chambers, b) in the forward drift chambers. Azimuth difference between a) and b) $\Delta\phi < 30$ mrad.	94
3. Track/cluster match	Track impact to cluster centroid distance $\Delta < 1$ cell size. Compare the energy distribution observed in the 3×3 cells centred on the impact cell to that expected for an electron incident along the track. Light sharing between phototubes of impact cell and hadronic leakage are taken into account. Require that the probability P that the energy distribution for an electron is farther from the mean than that observed is larger than 10^{-4} .	95	Shower position calculated from light sharing between impact cell phototubes must be consistent with the track impact point to within $\Delta x \leq 100$ mm	98
4. Presence of a pre-shower counter signal	Measured in C5 as a coincidence in space of charge clusters measured on the anode wires and the inner and outer cathode strips. Its charge q_s must exceed 3 minimum ionizing particles.	$> 95^a)$	Measured in MTPC in both coordinate planes. Its charge $Q(\text{MTPC})$ must exceed 6 minimum ionizing particles	93
5. Track/pre-shower position match	Within $d \leq 7$ mm measured on the C5 surface	95	Within $\Delta x, \Delta y < 20$ mm in each direction	99
6. Momentum measurement			Momentum p measured in spectrometer and energy E measured in calorimeter must satisfy $ p^{-1} - E^{-1} / \sigma(p^{-1} - E^{-1}) < 4$	$95^b)$

a) This does not include a 5% inefficiency due to a $\Delta\phi = 18^\circ$ azimuthal region in which C5 was not operational because of electrical breakdown.

b) This takes radiative corrections into account.

*) The strict cuts are defined by: $H < 0.023 + 0.034 \ln E_{cl}$ (see cut 1), $P > 0.01$ (see cut 3), $d \leq 5.5$ mm (see cut 5).

Table 7
Properties of the UA1 dielectron events

Run, event	Properties of the individual electrons					Lepton-pair properties			General event properties			
	CD measurement			Calorimetry		Mass ^{c)} (GeV/c ²)	p _T (GeV/c)	x _F ^{d)}	E _{tot} (GeV)	Σ E _T (GeV)	Missing E _T (GeV)	Charged tracks
	p (GeV)	Δp ^{a)} (GeV)	Q	e.m. energy (GeV)	Had. energy (GeV)							
A 7433 1001	33	+9 -6	+	44	0.0	91±5	2.9±0.9	0.02±0.01	274	82	2.1±3.6	27
	63	+23 -13	-	48	0.0							
B 7434 746	27	+19 -8	+	42	0.1	97±5	7.9±1.2	0.39±0.01	494	149	9.3±5.0	67
	93	+66 -28	-	102	0.0							
C 6059 1010	32	+11 -6	+	61	0.0	98±5	8.0±1.5	0.17±0.01	412	143	3.3±4.8	38
	9	+1 -1	-	48	0.0							
D 7739 1279	b)	b)	b)	51	0.0	95±5	8.4±1.4	0.17±0.01	493	157	0.8±5.0	54
	50	+50 -17	-	55	0.1							

a) ±1σ including systematic errors.

b) Unmeasured owing to large dip angle.

c) These errors have been scaled up arbitrarily to 5 GeV to represent the present level of uncertainty in the overall calibration of the e.m. calorimeter.

d) x_F is defined as the longitudinal momentum of the dilepton divided by beam energy.

Table 8
Properties of the UA2 dielectron events

Event	Properties of individual electrons			Lepton-pair properties	
	Region ^{a)}	Calorimetry		Mass (GeV/c ²)	p _T (GeV/c)
		e.m. energy (GeV)	Hadr. energy (%)		
A	C	49.7±1.6	6.5	91.6±1.9	0.4±1.4
	F	71.5±1.8	0.4		
B	C	48.4±1.7	5.0	94.6±2.4	12.7±1.8
	C	72.1±2.6	12.9		
C	C	11.5±0.7	5.0	90.6±1.9 ^{b)}	1.4±0.6
	F	69.9±1.8	0.7		
D	C	51.3±1.6	8.1	91.9±3.2	2.7±1.7
	F	104.5±6.4	0.4		
E	C	38.6±1.4	6.2	95.5±2.3	6.2±1.1
	F	60.0±2.0	0.3		
F	C	55.6±2.1	10.8	92.0±2.3	5.5±2.0
	C	46.6±1.6	8.7		
G	C	45.1±1.5	8.9	83.4±1.9 ^{c)}	10.4±1.8
	C	47.6±1.5	3.1		
H	C	46.3±1.6	10.6	88.0±2.1 ^{c)}	6.7±0.9
	C	42.1±1.4	4.4		

a) C = central, F = forward.

b) This is the (e⁺e⁻γ) mass.

c) Ignoring additional energy measured in neighbouring cells. Its inclusion results in a mass increase of ~ 3 GeV/c².

Table 9
Properties of the UA1 dimuon events

Run, event	Properties of the individual muons			Lepton-pair properties			General event properties			
	$p^a)$ (GeV/c)	Δp (GeV/c)	Q	Mass (GeV/c ²)	p_T (GeV/c)	x_F	E_{tot} (GeV)	$\Sigma E_T $ (GeV)	Missing E_T (GeV)	Charged tracks
E 6219 947	34.4	$+6$ -5	+	80.1^{+13}_{-10}	18.1^{+9}_{-7}	0.01	383	83.1	11.4 ± 7.8	50
	47.8	$+11$ -7	-							
F 6600 222	55.8	$+8$ -6	+	86.5^{+9}_{-8}	21.2^{+6}_{-5}	0.15	167	47.3	16.6 ± 5.2	26
	40.1	$+5$ -4	-							
G 6636 509	50.6 ^{b)}	$+6$ -5	+	$88.4^{+46^{c)}}_{-15}$	11.7^{+43}_{-14}	0.33	219	26.1	5.7 ± 5.9	30
	42.2 ^{b)}	$+44$ -14	-							
H 8523 831	61.8	$+16$ -11	+	95.8^{+24}_{-11}	2.5^{+16}_{-11}	0.20	270	27.7	1.3 ± 6.6	22
	45.7	$+12$ -8	-							
I 7428 1110	40.6	$+36$ -13	+	87.6^{+61}_{-28}	6^{+48}_{-19}	0.03	303	34.0	5.6 ± 7.0	32
	47.5	$+34$ -14	-							

a) This is the average momentum obtained from the three methods.

b) Momenta obtained without E_T balance.

c) This is the $(\mu^+\mu^-\gamma)$ mass.

Table 10

 $W^\pm$ and Z^0 parameters from the UA1 and UA2 experiments

	UA1	UA2
$N(W \rightarrow e\nu)$	52 ^{a)}	37 ^{b)}
$m_W (\text{GeV}/c^2)$	$80.9 \pm 1.5 \pm 2.4$	$83.1 \pm 1.9 \pm 1.3$
$\Gamma_W (90\% \text{ CL})$	$\leq 7 \text{ GeV}$	—
$(\sigma_B) (\text{nb})$	$0.53 \pm 0.08 \pm 0.09$	$0.53 \pm 0.10 \pm 0.10$
$N(W \rightarrow \mu\nu)$	14	—
$m_W (\text{GeV}/c^2)$	81.0^{+6}_{-7}	—
$(\sigma_B) (\text{nb})$	$0.67 \pm 0.17 \pm 0.15$	—
$N(Z^0 \rightarrow e^+e^-)$	$3 + 1^c)$	$7 + 1^c)$
$m_{Z^0} (\text{GeV}/c^2)$	$95.6 \pm 1.4 \pm 2.9$	$92.7 \pm 1.7 \pm 1.4$
$\Gamma_{Z^0} (90\% \text{ CL})$	$\leq 8.5 \text{ GeV}$	$\leq 6.5 \text{ GeV}$
$(\sigma_B) (\text{nb})$	$0.05 \pm 0.02 \pm 0.009$	$0.11 \pm 0.04 \pm 0.02$
$N(Z^0 \rightarrow \mu^+\mu^-)$	$4 + 1^c)$	—
$m_{Z^0} (\text{GeV}/c^2)$	85.6 ± 6.3	—
$(\sigma_B) (\text{nb})$	$0.105 \pm 0.05 \pm 0.15$	—
$\sin^2 \theta_w = 38.5/m_W$	0.226 ± 0.015	$0.216 \pm 0.010 \pm 0.007$
$\rho = [m_W/m_Z \cos \theta_w]^2$	0.968 ± 0.045	1.02 ± 0.06

a) $p_T^e > 15 \text{ GeV}/c$ b) $p_T^e > 25 \text{ GeV}/c$ c) $Z^0 \rightarrow \ell^+\ell^-\gamma (E_\gamma > 20 \text{ GeV})$

Table 11

Properties of the $\ell^+\ell^-\gamma$ events

	$e^+e^-\gamma$ (UA1)	$e^+e^-\gamma$ (UA2) ^{a)}	$\mu^+\mu^-\gamma$ (UA1)
E_γ (GeV)	38.8 ± 1.5	24.4 ± 1.0	28.3 ± 3
E_{ℓ^+} (GeV)	61.0 ± 1.2	69.9 ± 1.8	$50.6^{+5.8}_{-4.7}$
E_{ℓ^-} (GeV)	9 ± 1	11.5 ± 0.7	$42.2^{+44.0}_{-14.3}$
$\Delta\alpha(\ell^+, \gamma)(^\circ)^{b)}$	132.0 ± 4.0	129.9	7.9
$\Delta\alpha(\ell^-, \gamma)(^\circ)^{b)}$	14.4 ± 4.0	31.8	99.0
$m(\ell^+\ell^-)$ (GeV/c ²)	42.7 ± 2.4	50.4 ± 1.7	$70.9^{+37.2}_{-12.4}$
$m(\ell^+\ell^-\gamma)$ (GeV/c ²)	98.7 ± 5.0	90.6 ± 1.9	$88.4^{+46.1}_{-15.2}$
$m(\ell^+\gamma)$ (GeV/c ²)	88.8 ± 2.5	74.7 ± 1.8	5.0 ± 0.4
$m(\ell^-\gamma)$ (GeV/c ²)	4.6 ± 1.0	9.1 ± 0.3	$52.5^{+27.5}_{-9.3}$

a) The e^+ is identified since it goes into the forward region; the charge of ℓ^- is unknown.b) $\Delta\alpha$ is the angular difference in space.

Figure captions

- Fig. 1: Four-fermion interaction for neutron β -decay.
- Fig. 2: Neutron β -decay by W-exchange.
- Fig. 3: a) Feynman diagram for the reaction $\nu + \bar{\nu} \rightarrow W^+ + W^-$, leading to divergences.
b) Feynman diagram for the reaction $\nu + \bar{\nu} \rightarrow W^+ + W^-$ by Z^0 -exchange, compensating the divergences of diagram (a).
- Fig. 4: Neutral leptonic reaction $\bar{\nu}_\mu^{(-)} + e^+ \rightarrow \bar{\nu}_\mu^{(-)} + e^-$ mediated by a Z^0 -boson.
- Fig. 5: a) The original Drell–Yan mechanism.
b) The Drell–Yan process for the production of the weak bosons.
- Fig. 6: a) Differential luminosity $\tau(dL/d\tau)$ as a function of $\sqrt{\tau}$ for $u\bar{u}$ collisions in $p\bar{p}$ collisions. The upper curve (for $\sqrt{\tau} > 0.1$) is calculated without scaling violation, while the lower one includes scaling violations [from Ref. 49].
b) Differential luminosity for $u\bar{u}$ collisions in $p\bar{p}$ (upper) and pp (lower) collisions for $\sqrt{s} = 540$ GeV. Scaling violations are included [from Ref. 49].
- Fig. 7: Predicted production cross-sections for $W^\pm$ and Z^0 in pp and $p\bar{p}$ collisions. The calculated cross-sections include scaling violations (from Ref. 58).
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- Fig. 37: a) Fraction of energy deposited in the first sample of the gondola for all isolated electrons.
b) Fraction of energy deposited in the fourth sample of the gondola for all isolated electrons.
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- Fig. 38: Normalized deviation between the track hit as computed from the central detector and calorimetry centroids. The deviations have been measured in test beam runs
a) in azimuth around the beam axis,
b) in position, along the beam direction.
Arrows show the cuts imposed on the isolated electron selection.
- Fig. 39: Magnetic deflection in $1/p$ units compared to the inverse of the energy deposited in the gondolas for all isolated electron candidates with missing energy $E_T > 15 \text{ GeV}$.
- Fig. 40: a) The digitizations from the central detector for the tracks in one of the events which has an identified, isolated, well-measured high- p_T electron. The arrow points to the electron track.
b) The energy deposited in the cells of the central calorimetry and the equivalent plot for track momenta in the central detector for the event shown in (a). The top diagram shows the electromagnetic cells, the middle one shows the central detector tracks, and the bottom plot, with a very much increased sensitivity, shows the energy in the hadron

calorimeter. The plots reveal no hadronic energy behind the electron and no jet structure.

- Fig. 41: Transverse energy distributions of the electrons for $W \rightarrow e\nu$ events. The curve represents the result of a simple Monte Carlo simulation of a two-body decay, including the resolutions on electron and neutrino energies as well as the measured transverse momentum of the $(e\nu)$ system.
- Fig. 42: Transverse energy distribution of the neutrinos for $W \rightarrow e\nu$ events. The curve results from the same Monte Carlo simulation as in Fig. 41.
- Fig. 43: The missing transverse energy (E_ν) plotted vectorially against the electron direction for the $W \rightarrow e\nu$ events.
- Fig. 44: a) The component of the missing energy parallel to the electron momentum plotted against the electron energy for the $W \rightarrow e\nu$ events.
b) The component of the missing transverse energy (E_ν) perpendicular to the electron momentum, plotted against the electron transverse energy for $W \rightarrow e\nu$ events.
- Fig. 45: The transverse-momentum distribution of the $W \rightarrow e\nu$ events using the electron and missing transverse-energy vectors. The highest $p_T^{(W)}$ events have a visible jet (shown in black in the figure). The data are compared with the theoretical predictions of Halzen et al. [111] for W-production (a) without $O(\alpha_s)$ and (b) with QCD smearing; and predictions by (c) Aurenche and Kinnunen [112], and (d) Nakamura et al. [113].
- Fig. 46: a) The distribution of the transverse mass derived from the measured electron and neutrino vectors. The two curves show the results of a fit to the hypotheses $W \rightarrow e\nu$ and $X \rightarrow e\nu\nu$.
b) The enhanced electron-neutrino transverse-mass distribution (see text). The two curves show the results of a fit to the hypotheses $W \rightarrow e\nu$ and $X \rightarrow e\nu\nu$.
- Fig. 47: The fractional beam energy x_W carried by the W. The curve is the prediction obtained by assuming the W has been produced by $q\bar{q}$ fusion [58]. Note that in general there are two kinematic solutions for x_W (see text), which are resolved in 70% of the events by consideration of the energy flow in the rest of the event. Where this ambiguity has been resolved the preferred kinematic solution has been the one with the lowest x_W . In the 30% of the events where the ambiguity is not resolved the lowest x_W solution has therefore been chosen.
- Fig. 48: a) The x-distribution of the proton quarks producing the W by $q\bar{q}$ fusion. The curve is the prediction assuming $q\bar{q}$ fusion [58].
b) The same as (a) for the antiproton quarks.
- Fig. 49: a) The x-distribution of the u ($\bar{u}$) quarks in the proton (antiproton).
b) The same as (a) but for d ($\bar{d}$) quarks in the proton (antiproton).
- Fig. 50: Three-dimensional representation of $W^\pm$ population as a function of electron angle and transverse momentum. The lines represent (theoretically predicted) levels of equal population, with an arbitrary normalization. The circles represent the observed W^- events, the crosses the observed W^+ events.
- Fig. 51: Angular distribution of the electron emission angle θ^* in the rest frame of the W after correction for experimental acceptance. Only those events are included for which both the charge of the electron and the angle θ^* are unambiguously determined. The prediction of a $V \pm A$ theory is shown for comparison.

- Fig. 52: Electron transverse-momentum distribution for $W \rightarrow e\nu$ events. The full line is the background estimate and the dashed line is the expectation from $W \rightarrow e\nu$.
- Fig. 53: Electron transverse-momentum distribution for $W \rightarrow e\nu$ events after applying strict electron identification criteria.
- Fig. 54: Corrected transverse-momentum distribution of the 40 $W \rightarrow e\nu$ events together with the QCD expectation.
- Fig. 55: a) The component of the neutrino momentum parallel to the electron plotted against the electron transverse momentum for the $W \rightarrow e\nu$ events.
b) The component of the neutrino momentum perpendicular to the electron plotted against the electron transverse momentum for the $W \rightarrow e\nu$ events.
- Fig. 56: Distribution of the $W \rightarrow e\nu$ events with known charge for the electron in the $(1/E, 1/p)$ plane.
- Fig. 57: A graphical display of a $W^- \rightarrow \mu^+\nu$ event. The vertical arrow shows the trajectory of the 25 GeV/c μ^+ up to the muon chamber, while the other arrow shows the transverse direction of the neutrino. The curved lines from the vertex are the charged tracks seen by the central detector, and the petals and boxes illustrate the electromagnetic and hadronic energy depositions. An expanded view of a muon module is shown as an insert.
- Fig. 58: Two-dimensional plot of Q/p_μ versus Q/p_{CD} for the $W \rightarrow \mu\nu$ events, where p_μ and p_{CD} are the momenta measured in the muon chambers and central detector respectively, and Q is the charge of the muon. Two events with tracks in the bottom chambers are not shown.
- Fig. 59: Transverse energy of the neutrino parallel to the muon versus transverse momentum of the muon. Since the two quantities are correlated, error bars are shown for the difference and the sum. The difference is the transverse energy of the W parallel to the muon, which is measured in the calorimetry and is therefore not correlated with the transverse momentum of the muon. For the errors in the sum only two error bars are shown for typical events. The filled circles correspond to the final sample of 14 W events, and the open circles to the 4 events with $p_T^\mu < 15$ GeV/c.
- Fig. 60: The solid curve is an ideogram of the transverse momentum distribution of the muons in the final sample of 14 $W \rightarrow \mu\nu$ events. The dashed curve is a Monte Carlo prediction, based on the W production spectra measured in $W \rightarrow e\nu$ decays and a W mass of 80.9 GeV/c² smeared with errors.
- Fig. 61: The transverse-momentum distribution of the W , derived from the energy imbalance measured in the calorimetry. The corresponding distribution from the $W \rightarrow e\nu$ data is shown for comparison.
- Fig. 62: Invariant mass distribution (uncorrected) of two electromagnetic clusters:
- Each electromagnetic cluster has $E_T > 25$ GeV;
 - As in (a) but in addition a track with transverse momentum $p_T > 7$ GeV/c or compatible with $p_T \geq 15$ GeV/c must point to one of the clusters.
 - The cluster hit by the charged track must be isolated. In addition, a small energy deposition in the hadron calorimeter immediately behind (< 0.6 GeV) ensures the electron signature.
 - The second cluster also contains a high- p_T track, is isolated, and shows an electromagnetic signature.

- Fig. 63: a) Event display. All reconstructed vertex-associated tracks and all calorimeter hits are displayed.
b) The same, but thresholds are raised to $p_T > 2 \text{ GeV}/c$ for charged tracks and $E_T > 2 \text{ GeV}$ for calorimeter hits. We remark that only the electron pair survives these mild cuts.
- Fig. 64: Electromagnetic energy depositions at angles $> 5^\circ$ with respect to the beam direction for the four electron pairs.
- Fig. 65: Normalized deviation between the track hit from the central detector as compared to the calorimeter centroids for a) the position along the beam direction, (the curve is the expected distribution from test-beam data) and b) the azimuth around the beam axis.
- Fig. 66: Magnetic deflection in $1/p$ units compared to the inverse of the energy deposited in the electromagnetic calorimeters. Ideally, all electrons should lie on the $1/E = 1/p$ line.
- Fig. 67: Invariant mass distributions:
a) of the 24 pairs which pass cut 1 of Table 6.
b) of the eight of these 24 pairs for which all cuts of Table 6 are satisfied by at least one electron; the three events in which both electrons pass all cuts are cross-hatched.
- Fig. 68: a) Longitudinal view of event A in the plane containing the central electrons. In each of the four proportional chambers of the vertex detector (PROP 1 to 4) and in the pre-shower chamber C5 located behind the tungsten converter, signals are indicated whenever a coincidence in space was observed between the anode wire and the inner and outer cathode strips. The measurements from the two drift chambers (DRIFT 1 and 2) are indicated as crosses with sizes corresponding to the uncertainty on the charge-division measurement. Energies measured in the electromagnetic cells facing the electron track are indicated, when non zero.
b) The transverse-energy distribution in the cells for event A in the (θ, ϕ) plane.
- Fig. 69: a) The transverse-energy distribution in the cells for event C in the (θ, ϕ) plane. Electron e1 and the photon (γ) are observed in the central region, electron e2 in the forward region.
b) The transverse view of event C. Signals from the proportional and drift chambers are indicated as dots. Electron e1 and the photon (γ) are observed in the central region and associated with C5 signals (indicated by heavy lines proportional to pulse height). Electron e2 is observed in the forward region (not covered by C5).
- Fig. 70: Invariant dimuon mass.
- Fig. 71: Display for a high-invariant-mass muon-pair event: (a) without cuts, and (b) with $p_T > 1 \text{ GeV}$ thresholds for tracks and $E_T > 0.5 \text{ GeV}$ for calorimeter hits.
- Fig. 72: Normalized energy losses in calorimeter cells traversed by the muon tracks.
- Fig. 73: Arrows show residuals in angle and position for the muon tracks. Distributions come from cosmic-ray calibration with $p > 50 \text{ GeV}/c$.
- Fig. 74: Invariant masses of lepton pairs from the UA1 and UA2 experiments.
- Fig. 75: The transverse momentum distribution of the Z^0 derived from the $Z^0 \rightarrow e^+e^-$ and $Z^0 \rightarrow \mu^+\mu^-$ events of the UA1 and UA2 experiments. The theoretical prediction is shown for comparison.
- Fig. 76: Comparison between the standard model expectations [126] and the experimental results combined from the UA1 and UA2 experiments.

Fig. 77: One-sided confidence levels in the plane Γ_z and $(\sigma B)_z/(\sigma B)_w$, and determination of the number of neutrino types.

Fig. 78: Upper limit (90% CL) for R , the fraction of leptonic W decays accompanied by a photon with transverse energy in excess of 10 GeV. The electron samples have been used to extract the limits on R , which are model-dependent. Results are shown for the following two models:

- a) the process $W \rightarrow e^* \nu_e (e^* \rightarrow e \gamma)$ with $m_W = 80.9 \text{ GeV}/c^2$, and
- b) $\bar{W} \rightarrow W \gamma (W \rightarrow e \nu_e)$, where one of the boson masses has been taken to be $80.9 \text{ GeV}/c^2$, and the shaded band shows the dependence of the limit on the boson spin assignments (0 or 1).

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5

a)

b)

Fig. 6

Fig. 7

Fig. 8

a) $W^{\pm}$ production

b) $W^{\pm}$ leptonic decay:

Fig. 9

Fig. 10

Fig. 11

Fig. 12

Fig. 13

Fig. 14

a)

b)

Fig. 16

Fig. 15

a)

b)

Fig. 20

Fig. 19

Fig. 21

Fig. 22

Fig. 23

Fig. 24

Fig. 25

Fig. 26

TRANSVERSE VIEW OF VERTEX DETECTOR

Fig. 27

Fig. 28

Fig. 29

Fig. 30

Fig. 31

Fig. 32

Fig. 33

Fig. 34

Fig. 35

Fig. 36

Fig. 37

(Matching in azimuth)

Fig. 38

Fig. 39

RUN 4017 Evt 838 b)

a)

EVENT 4017. 838.

Fig. 40

Fig. 41

Fig. 42

Fig. 43

Fig. 44

Fig. 45

Fig. 46

Fig. 47

Fig. 48

Fig. 49

Fig. 50

Fig. 51

Fig. 52

Fig. 53

Fig. 54

Fig. 55

Fig. 56

EVENT 7163. 238. THR 0.50 PTCO 0.00 ETTH 0.0

Fig. 57

Fig. 58

Fig. 59

Fig. 60

Fig.61

Fig.62

Fig. 63

Fig. 64

Fig. 65

Fig. 66

Fig. 67

Fig. 68

Fig. 69

Fig. 70

EVENT 6600. 222.

EVENT 6600. 222.

Fig. 71

Fig. 72

Fig. 73

Fig. 74

Fig. 75

Fig. 76

Fig. 77

Fig. 78