FTUV/91-53
IFIC/91-52

ELECTRO-WEAK INTERACTION IN MUONIC ATOMS [1]

J. BERNABEU

Departament de Física Teòrica,
Universitat de València, e
IFIC, Centre Mixte Univ. València-CSIC,
46100 Burjassot, València
SPAIN

Abstract

The parity non-conserving effective neutral current interaction between charged leptons and nucleons is studied in its implications for atomic physics. Present results on heavy electronic atoms are discussed within the standard electro-weak theory and beyond. The new features provided by muonic atoms open the way to the nuclear-spin- dependent parity non-conserving effects. Different observables proposed to study these effects in muonic atoms are reviewed.

[1] Dedicated to V.W.Hughes and G. zu. Putlitz. Invited Lecture at the Workshop "The Future of Muon Physics", Heidelberg, May 1991

1 INTRODUCTION

The standard theory of electro-weak interactions has been confirmed by high precision experiments [1] at the CERN electron-positron collider LEP. The measurements of the Z-mass, its width, the leptonic widths, the hadronic width, the forward-backward asymmetries, the longitudinal polarization of the τ-lepton,..., have given a beautiful test of the standard theory and a precise determination of its parameters. Apart from their intrinsic interest, parity non-conserving (PNC) observables in Atomic Physics are complementary to LEP. In this paper I would like to discuss the physics content of the PNC effective neutral current interaction between charged leptons and quarks, its implications for atomic physic and the role played by PNC observables in muonic atoms.

The effective Lagrangian for the PNC lepton-quark interaction is written as

$$L_{PNC} = \frac{G_F}{\sqrt{2}} \sum_i \left[C_{1i} \bar{l}\gamma^\mu \gamma_5 l \bar{q}_i \gamma_\mu q_i + C_{2i} \bar{l}\gamma^\mu l \bar{q}_i \gamma_\mu \gamma_5 q_i \right] \tag{1}$$

where the index $i = u, d, s...$, runs over the different flavours of quarks. The Standard Model at tree level gives [2] the following couplings:

$$C_{1u} = -\frac{1}{2} + \frac{4}{3} sin^2\Theta_w$$

$$C_{1d} = \frac{1}{2} - \frac{2}{3} sin^2\Theta_w \tag{2}$$

for the axial current of leptons times the vector current of quarks, and

$$C_{2u} = -C_{2d} = -\frac{1}{2} + 2 sin^2\Theta_w \tag{3}$$

for the vector current of leptons times the axial current of quarks. Using the weak isospin symmetry for the different families of quarks, we have $C_{1s} = C_{1d}$ and $C_{2s} = C_{2d}$.

The matrix elements of the hadronic current for the nucleon are written for $Q^2 = O$ as

$$< p|\bar{q}_i \gamma_\mu q_i|p >= G_V^{(i)} \bar{p}\gamma_\mu p$$

$$< p|\bar{q}_i\gamma_\mu\gamma_5 q_i|p >= G_A^{(i)}\bar{p}\gamma_\mu\gamma_5 p \tag{4}$$

The conserved vector current and its coherent character, with the vector charge equal to the quark-number, determine the couplings $G_V^{(u)} = 2, G_V^{(d)} = 1, G_V^{(s)} = 0$, for the proton.

In terms of definite $U(3)$ flavour transformation properties, one can introduce the following combination of couplings

$$G_A^{(3)} = G_A^{(u)} - G_A^{(d)}$$

$$G_A^{(8)} = G_A^{(u)} + G_A^{(d)} - 2G_A^{(s)} \tag{5}$$

$$G_A^{(0)} = G_A^{(u)} + G_A^{(d)} + G_A^{(s)}$$

The charged weak currents transform, following Cabibbo theory for the light quarks, as an octet under flavour SU(3). Then the two form factors $G_A^{(3)}$ and $G_A^{(8)}$ can be expressed through the amplitudes F and D known from semi-leptonic decays of baryons

$$G_A^{(3)} = F + D = 1.254 \pm 0.006 \ ; \ G_A^{(8)} = 3F - D = 0.68 \pm 0.04 \tag{6}$$

The EMC measurement [3] of the polarization-dependent structure function of the proton determines an independent combination of $G_A^{(3)}$, $G_A^{(8)}$ and the singlet $G_A^{(o)}$. One obtains

$$G_A^{(o)} = 0.12 \pm 0.17 \tag{7}$$

It is remarkable that the singlet current coupling $G_A^{(o)}$ is compatible with zero, which seems to be in contradiction with naive expectations from models with constituent u , d quarks for the proton. In such models it would mean that the total helicity carried by all quarks (and antiquarks) in a polarized proton is small. For our considerations we do not explicitly need constituent quark models of the nucleon, since we work instead with the effective nucleonic currents of Eq.(4). The above results lead to

$$G_A^{(s)} = -0.19 \pm 0.06 \tag{8}$$

This important conclusion would be confirmed in the SMC-experiment at CERN, conducted by Prof. V.Hughes.

Using strong isospin as a symmetry of the strong interactions in the limit in which $(m_d - m_u)$ is small, we have that for the neutron the following substitutions have to be made

$$p \to n \Rightarrow \begin{array}{c} G_{V,A}^{(u)} \rightleftharpoons G_{V,A}^{(d)} \\ G_{V,A}^{(s)} \to G_{V,A}^{(s)} \end{array} \tag{9}$$

where $G_{V,A}^{(i)}$ are the form factors defined in Eq.(4).

2 PARITY NON-CONSERVATION IN ATOMS

The study of PNC observables in Atomic Physics provides a test of the fundamental symmetries of Nature and, specifically, of the standard electro-weak theory, the neutral current lepton-quark couplings and the corresponding electroweak radiative corrections.

Parity Violation effects have been observed in heavy electronic atoms [4]. They are characterized by different experimental methods, atomic elements and transitions observed. All of them use the enhancement, pointed out by Bouchiat and Bouchiat, of the neutral current mixing parameter in heavy atoms. The piece of the effective electron-quark parity violating interaction corresponding to the axial coupling of the electron times the vector coupling of quarks becomes coherent in the nucleus, so the expectation value is proportional to the number of quarks. The corresponding weak neutral charge induced by the PNC Z-exchange diagram, which interferes with γ-exchange as shown in Figure 1, is given by

$$Q_w = -2\{(2Z + N)C_{1u} + (Z + 2N)C_{1d}\} \tag{10}$$

where (Z,N) are the (proton, neutron) numbers in the nucleus and $C_{1u,d}$ are the couplings of Eq.(2).

The PNC interaction (with time-reversal symmetry) induces pure-imaginary atomic dipole moments that are observable in atomic transitions. There are two important

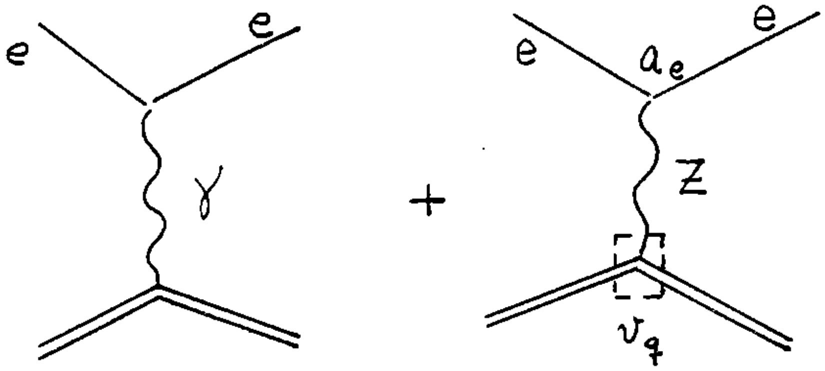

Figure 1

classes of such experiments, each involving the interference between an ordinary (but relatively weak, in order to enhance the effect) electromagnetic transition and the PNC-induced E1 transition. In the first (M1) class of experiments, the induced E1 transition is coupled into an M1 transition, producing circularly polarized light or a rotation of plane polarized light. The quantity which is measured by the experiments searching for optical rotation of light in atomic vapour is $R = Im(E_1^{PNC}/M_1)$, which factorizes Q_w from the atomic structure dependent inputs. In Table 1 the experimental values of the quantity R for bismuth (two transition lines), lead and thallium are collected.

Experiment	*Element(line)*	*R*	
Novosibirsk	$Bi(648)$	(-20.2 ± 2.7)	$.10^{-8}$
Oxford	$Bi(648)$	(-9.3 ± 1.5)	$.10^{-8}$
Moscow	$Bi(648)$	(-7.8 ± 1.8)	$.10^{-8}$
Seattle(1981)	$Bi(876)$	(-10.4 ± 1.7)	$.10^{-8}$
Seattle(1983)	$Pb(1280)$	(-9.9 ± 2.5)	$.10^{-8}$
Berkeley(1981)	Tl	(-14.0 ± 3.7)	$.10^{-4}$

Table 1

The difference in the order of magnitude of the effect for thallium is due to the choice of a forbidden magnetic transition in this case.

In the second (Stark) class of experiments, a static electric field $\vec{E}_s$ induces an E1 transition between two states, which interferes with the PNC-induced E1 transition between the same two states. The corresponding transition rate depends then on the relative directions of $\vec{E}_s$ and the polarization vectors of the absorbed and reemitted radiation. The measured quantity is expressed in terms of $Im(E_1^{PNC}/\beta)$, where β is the vector polarizability. The experimental values obtained are shown in Table 2:

Experiment	*Element*	$Im(E_1^{PNC}/\beta)(mV/cm)$
Berkeley(1981)	*Tl*	-1.80 ± 0.48
Berkeley(1981)	*Tl*	-1.73 ± 0.27
Paris(1982)	*Cs*	$+1.33 \pm 0.25$
Paris(1984)	*Cs*	-1.75 ± 0.27
Boulder(1985)	*Cs*	-1.65 ± 0.13
Boulder(1988)	*Cs*	-1.576 ± 0.034

Table 2:

The improved measurement [5] of PNC effects in cesium by the Boulder (1988) experiment provides an experimental precision of 2 %. This allows the extraction of electroweak radiative correction contributions if the atomic structure calculations are accurate enough. A major theoretical effort for the $6S_{1/2} \rightarrow 7S_{1/2}$ transition in Cs has been presented in Ref. [6], with the claim of 1% theoretical uncertainty. The extracted value of Q_w from the Boulder (1988) experimental result of Table 2 and this atomic structure calculation is

$$Q_w^{exp}(Cs) = -71.04(1.58)[0.88] \tag{11}$$

where the first error is from the experiment and the second from the atomic theory.

When the result (11) is interpreted in the standard electroweak theory, the inclusion of the electroweak radiative corrections brings not only $sin^2\Theta_w$ but also m_t as parameters of the PNC-observable. In the modified minimal-subtraction scheme, one gets [6]

$$sin^2\hat{\Theta}_w(m_w) = \begin{array}{l} 0.2242(65)[36], m_t = 100GeV \\ 0.2215(65)[36], m_t = 200GeV \end{array} \tag{12}$$

This result is compared with the determination provided by the precise Z-mass measurement [7] in the LEP experiments, as shown in Fig.2. The dependence on m_t is almost identical for $Q_w(Cs)$ and for m_Z:
The LEP-value of $sin^2\hat{\Theta}_w$ from the leptonic decay widths of the Z is also included in Fig. 2 for comparison.

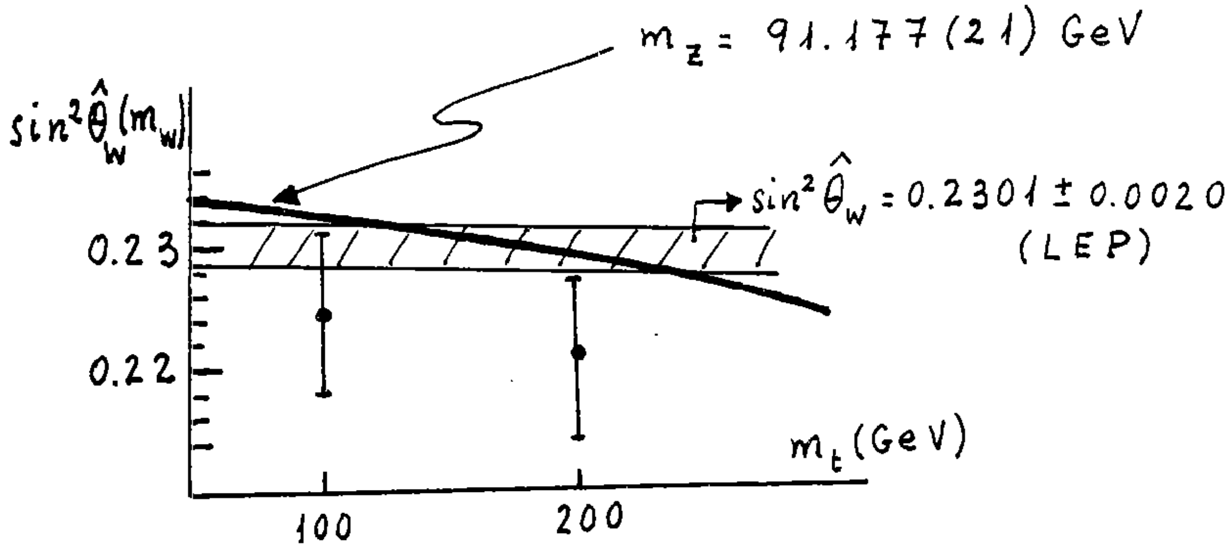

Figure 2

If, instead of using $sin^2\hat{\Theta}_w$ as independent parameter of the theory, one takes the precisely known Z-mass together with α (fine structure constant), G_F (Fermi coupling constant), m_t (top quark mass) and m_H (Higgs mass), the theoretical prediction of $Q_w(Cs)$ turns out to be almost independent of m_Z.

One gets

$$Q_w^{th}(Cs) = -73.10 \pm 0.13 \tag{13}$$

so that, independently of the unknown value of m_t, the difference

$$\delta Q_w(Cs) = Q_w^{exp} - Q_w^{th} = 2.1 \pm 1.8 \tag{14}$$

constitutes a test of the electroweak radiative corrections to the standard theory and a probe sensitive to new physics.

One example of the sensitivity of δQ_w to new physics, appearing through loop corrections, is given by technicoulor theories [8]. With the self-energy corrections evaluated in one-loop technifermion approximation, the value of Q_w receives a negative contribution, whose size is proportional both to the number of technicolours and of technidoublets.

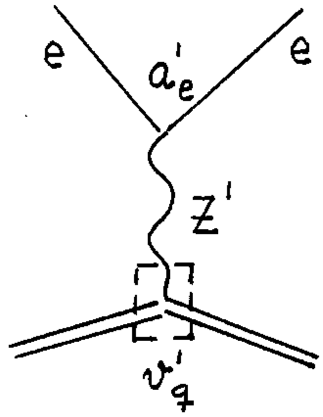

Figure 3

In models with an extra neutral vector boson Z', there is an additional tree-level contribution to Q_w coming from Z'-exchange, in the limit of small mixing with the standard Z. This is shown in Fig.3.

One has

$$\delta Q_w^{new} = 16\frac{m_Z^2}{m_{Z'}^2}[(2Z+N)a'_e v'_u + (Z+2N)a'_e v'_d] \tag{15}$$

where v'_f, a'_f are the vector and axial vector couplings of the fermion f to the vector boson Z'. These couplings depend [9] on the particular extended gauge theory, like models with extra $U(1)$ or left-right models. To be specific, a superstring-inspired model gives detailed predictions

$$v'_u = 0 \ , \ a'_e = v'_d = \frac{1}{4}\sqrt{\frac{5}{3}} sin\Theta_w \tag{16}$$

so that one would have

$$\delta Q_w^{new} = \frac{5}{3} sin^2\Theta_w (Z+2N) m_Z^2/m_{Z'}^2 \tag{17}$$

If the limit (14) for $\delta Q_w(Cs)$, whose value is compatible with zero, is used to bound $m_{Z'}$ from Eq.(17) one obtains a significant result $m_{Z'} \geq 400$ GeV.

3 NUCLEAR-SPIN-DEPENDENT EFFECTS

Among the virtues offered by the study of PNC-effects in muonic atoms, one encounters that they are practical tools to obtain a better access to all PNC-couplings of charged leptons to nucleons.

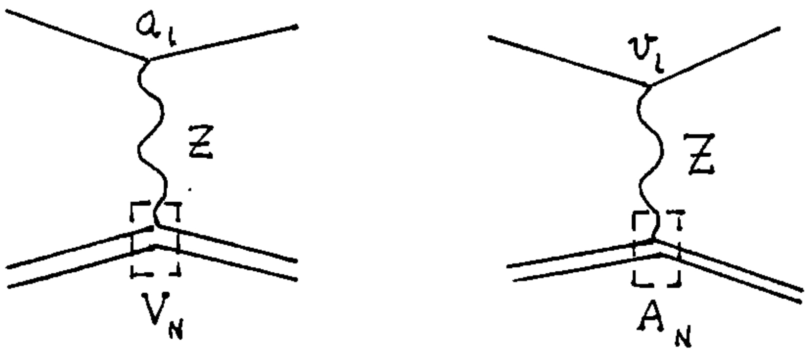

Figure 4

The effective Lagrangian for the PNC interaction of charged leptons with nucleons can be written as

$$L_{PNC} = \frac{G_F}{\sqrt{2}} \sum_{i=p,n} [C_{1i}\bar{l}\gamma^{\mu}\gamma_5 l \bar{N}_i\gamma_{\mu}N_i + C_{2i}\bar{l}\gamma^{\mu}l\bar{N}_i\gamma_{\mu}\gamma_5 N_i] \qquad (18)$$

associated with the two contributions from the Z-exchange diagram of Fig.4, respectively.

In terms of the neutral current lepton-quark couplings and the form factors of Section 1, one has for the nuclear-spin-independent couplings

$$C_{1p} = 2C_{1u} + C_{1d}$$

$$C_{1n} = C_{1u} + 2C_{1d} \qquad (19)$$

and their associated coherent action on the nucleus, as seen in Eq. (10). For the nuclear-spin-dependent terms one obtains [10]

$$C_{2p} = G_A^{(u)}C_{2u} + [G_A^{(d)} + G_A^{(s)}]C_{2d}$$

$$C_{2n} = G_A^{(d)}C_{2u} + [G_A^{(u)} + G_A^{(s)}]C_{2d} \qquad (20)$$

and they do not add coherently in the nucleus.

In order to single out the nuclear-spin-dependent PNC effect generated by the second term of Eq.(18) one has to measure PNC observables for different hyperfine

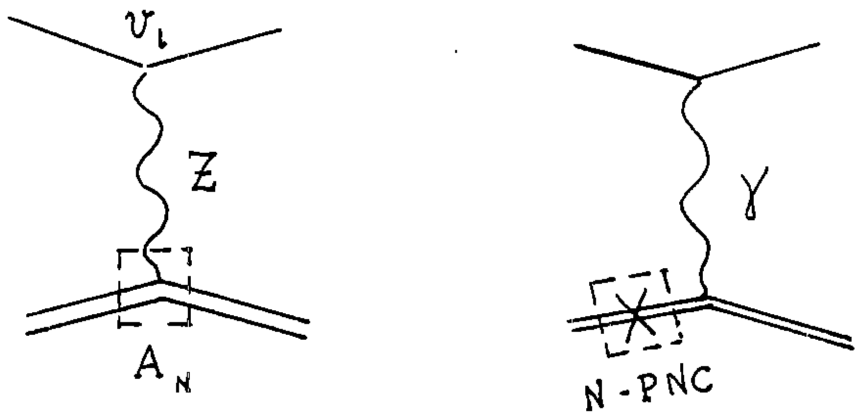

Figure 5: The weak neutral current and the nuclear anapole moment contributions to PNC effects

transitions in atoms. A first indication of this effect has been obtained by the Boulder (1988) experiment [5] in electronic Cs atoms. For different hyperfine levels, the measured quantity $Im(E_1^{PNC}/\beta)$ is different, as shown in Table 3:

	Lines	$Im(E_1^{PNC}/\beta)(mV/cm)$
$Boulder(1988), Cs$	$F=4 \rightarrow F'=3$	-1.639 ± 0.047
	$F=3 \rightarrow F'=4$	-1.513 ± 0.049

Table 3

The impressive precision reaching high levels of accuracy in the atomic experiments of the Stark-type can lead to significant nuclear- spin-dependent PNC effects. But the progress is difficult, because the experimental effect is comparatively very small. Furthemore, on the theoretical side the PNC-neutral axial current interaction contribution is contaminated by the effect of the nuclear anapole moment.

The two contributions which compete for the nuclear-spin-dependent PNC effect are given by the two diagrams of Fig.5.

The nuclear anapole moment describes [11] the effect of the parity violating nuclear forces on the nucleus electromagnetic current. This mechanism induces the same effective operator for the lepton-Nucleus amplitude as the neutral current interaction

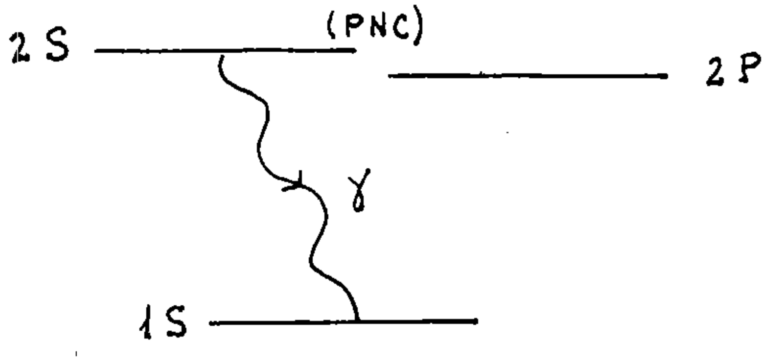

Figure 6

of the vector lepton current with the axial nuclear current. Although formally a higher order α-correction, the nuclear anapole moment contribution has the electromagnetic coherence in heavy nuclei whereas the neutral axial current effect has no coherent effect. Furthermore, the PNC lepton-nucleus interaction associated with the hadronic neutral axial current is suppressed due to the small vector lepton coupling. As a consequence, the nuclear-spin- dependent PNC effect is dominated by the electron interaction with the nuclear anapole moment in heavy atoms.

If one is interested in disentagling the neutral current effect, PNC-observables in light muonic atoms are thus of great value.

4 PARITY VIOLATION IN MUONIC ATOMS

The original proposal [12] to detect neutral currents in muonic atoms considered the transition between 2S and 1S states, as shown in Fig.6.

The 2S and 2P states are admixed due to the parity-odd piece of the neutral current interaction between muon and nucleons, as given by Eq.(18). The radiative transition between the 2S and 1S states is then a sum of magnetic and electric dipole amplitudes: M1 + (PNC)E1. The interference of the two amplitudes produces asymmetries in the distribution of the emitted radiation.

Among the interests in the search for PNC-effects in muonic atoms, one should point out the accessibility to the four effective couplings $C_{1p}, C_{1n}, C_{2p}, C_{2n}$ of Eq. (18), the different Q^2-value probed in muonic atoms, the complete control of the atomic physics structure (hydrogen-like), etc. Light muonic atoms offer the additional theoretical perspective of studying the nuclear-spin-dependent PNC-effects and, experimentally, they present important enhancements: 1) of the mixing parameter, due to the near degeneracy of the 2S and 2P levels, and 2) of the amplitude ratio $|E1/M1|$, which goes like $(\alpha Z)^{-3}$.

Interesting asymmetries are: 1) a net circular polarization of the radiation; 2) the angular asymmetry of the emitted radiation with respect to the polarization of the muon in the initial 2S- or final 1S- states, and 3) the directional correlation between the momenta of the emitted photon and the decay electron.

J. Missimer and L.M.Simons have proposed [13] the directional correlation as a suitable observable for the two isotopes of boron, ^{10}B and ^{11}B. Muonic boron is unique among the light muonic atoms in fulfilling the conditions necessary for the detection of the M1 transition:

1) Occuring in a gaseous state, the non-radiative 2S decays are suppressed;

2) The radiative lifetime of the 2S-state is much longer than that of other excited states, so it can be prepared by waiting for several ns after the muon stop, and

3) The M1 transition can be distinguished from the more probable two-photon (2E1) transition by using an X-ray absorption edge.

The two naturally occurring isotopes of boron have nonvanishing spins, so that measurements of the photon-electron directional correlation in the individual hyperfine components of the initial or final state of the radiative transition determines completely the neutral current couplings of the muon to nucleons. The energy levels of $\mu^{11}B$ relevant to parity-mixing experiments are shown in Fig.7.

The different asymmetries discussed above are interconnected in the following way. The partial rate of the hyperfine transition or the angular distribution of photons relative to the muon spin in the final 1S-state is given by

$$d\Gamma^{\gamma}_{F'\to F} \propto [d(F',F) + \vec{P}_{\mu}.\hat{k} \;\; n(F',F)]\frac{d\Omega}{4\pi} \tag{21}$$

where $\vec{P}_{\mu}$ is the μ-polarization and $\hat{k}$ the photon direction. We see that the forward-backward angular asymmetry coincides with the circular polarization of photons

$$P_{\gamma}(F',F) = \frac{n(F',F)}{d(F',F)} \tag{22}$$

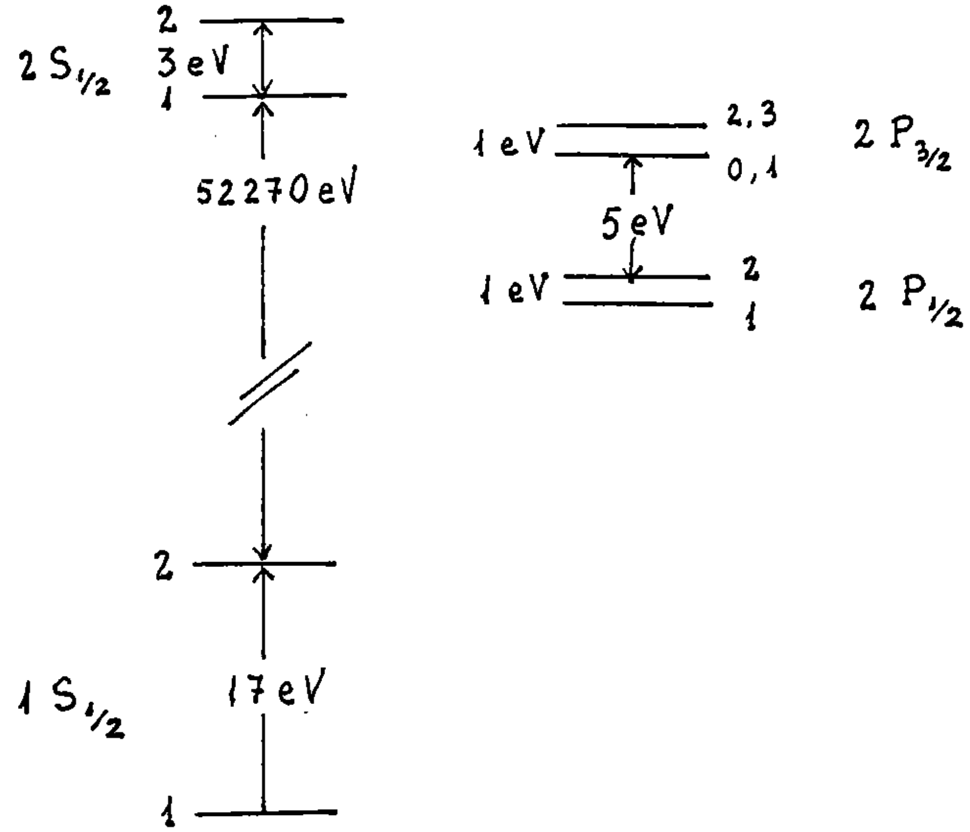

Figure 7

Since only two hyperfine states occur for each S-state, they can be labeled (+) for the $I + 1/2$ and (-) for the $I - 1/2$ components, where I is the nuclear spin. Thus, $P_\gamma(+,-)$ is the asymmetry in the transition between the $F' = I + 1/2$ component of the initial 2S-state and the $F = I - 1/2$ component of the final 1S-state.

The values of P_γ have been calculated for the $^{10}B -^{11} B$ isotopes [14] and for the $^3He -^4 He$ isotopes [15]. The results for boron are (in percent):

$$^{10}B(I=3) : P(+,+) = P(+,-) = 3.22[C_{1p} + C_{1n} + 0.20(C_{2p} + C_{2n})]$$

$$P(-,+) = P(-,-) = 3.52[C_{1p} + C_{1n} - 0.26(C_{2p} + C_{2n})]$$

$$^{11}B(I=3/2) : P(+,+) = P(+,-) = 3.34[C_{1p} + 1.2C_{1n} + 0.20C_{2p}] \tag{23}$$

$$P(-,+) = P(-,-) = 3.95[C_{1p} + 1.2C_{1n} - 0.34C_{2p}]$$

The equality of the asymmetries for identical initial states of the same atom follows from the absence of mixing in the final 1S-state. Inspection of the results (23) shows that the effects are large and that the measurement, if precise enough, of the different hyperfine transitions could determine the four interesting coupling constants. The relation of these effective couplings to the fundamental quark couplings is given in Eqs. (19) and (20), and it involves the flavour form factors $G_A^{(i)}$: the use of the nuclear-spin- dependent PNC effect in muonic atoms as a tool to determine the strangeness form factor $G_A^{(s)}$ has been emphasized in Ref. [10]. As an example, for the strong isoscalar ^{10}B nucleus the relevant couplings in the standard theory are

$$C_{1p} + C_{1n} = 2sin^2\Theta_w$$

$$C_{2p} + C_{2n} = [1 - 4sin^2\Theta_w]G_A^{(s)} \tag{24}$$

Alternatively, if the standard theory of electroweak interactions is not assumed, the PNC-effects in muonic atoms can be used to test extended gauge models, such as [16] extra Z-bosons, lepto-quarks, compositeness, etc. In this case, the aim is to extract the lepton-quark couplings from this experiment. Using the present results on $G_A^{(i)}$ discussed in Section 1, the information contained in Eq.(20) from C_{2p} and C_{2n} is

$$C_{2p} = (0.78 \pm 0.06)C_{2u} - (0.66 \pm 0.12)C_{2d}$$

$$C_{2n} = -(0.47 \pm 0.06)C_{2u} + (0.59 \pm 0.12)C_{2d} \tag{25}$$

which in both cases is almost proportional to the combination $C_{2u} - C_{2d}$.

In Ref.[13] the connection between P_γ of Eq.(22) and the directional correlation between the photon and the decay electron is shown:

$$W(\varepsilon, x) = [2\varepsilon^2(3 - 2\varepsilon) + 2\varepsilon^2(1 - 2\varepsilon)P_\gamma x]d\varepsilon\frac{dx}{2} \tag{26}$$

where ε is the fractional energy of the electron $\varepsilon \simeq \frac{E_e}{m_{\mu/2}}$, and $x = cos\Theta$ between the photon and electron momenta. The correlation manifests itself as an asymmetry A in the number of electrons emitted parallel and antiparallel to the photon direction.

The measurability is determined by the number of correlated events per second I occurring in given ranges of ε and x. The number of correlated events which must be measured to determine the asymmetry to a desired accuracy is inversely proportional to the figure of merit $I.A^2$. The maximun figure of merit for boron is estimated as max $(I.A^2) \simeq 1.10^{-11}$, so that 1.10^{13} μ-stops are required to measure the asymmetry P_γ to a relative accuracy of ten percent.

5 CONCLUSIONS

The nuclear-spin-independent PNC-effects in heavy electronic atoms are reaching high precision levels. Within the standard electroweak theory, the experimental results on Atomic Parity Violation constitute a very clean test of the electroweak radiative corrections. If interpreted with models beyond the standard theory, one gets a powerful exploration of additional Z-bosons, technicolour ideas, etc.

Parity Violation in Light Muonic Atoms offers a tool to separate out the nuclear-spin-dependent PNC-interaction. With appropriate selection of the nuclear isotopes one could separate the effective axial isovector $(C_{2p}-C_{2n})$ coupling of the nucleon from the effective axial isoscalar $(C_{2p}+C_{2n})$ coupling. They are related to the fundamental quark couplings and the nucleon structure through

$$C_{2p} - C_{2n} = (G_A^{(u)} - G_A^{(d)})(C_{2u} - C_{2d})$$

$$C_{2p} + C_{2n} = (G_A^{(u)} + G_A^{(d)})(C_{2u} + C_{2d}) + 2G_A^{(s)} C_{2d} \tag{27}$$

As seen, the accessibility to the $G_A^{(s)}$ form factor is of the highest interest for the understanding of the nucleon structure.

ACKNOWLEDGEMENTS

I would like to thank the organizers of the Workshop "The Future of Muon Physics" for the warm atmosphere that we had, which triggered lively discussions. This work is dedicated to Profs. V.W.Hughes and G. zu Putlitz, on the occasion of their anniversaries. Fruitful discussions with Prof. S. Bilenky are acknowledged. This research has been supported by CICYT under Grant AEN 90-0040.

REFERENCES

1. J.Steinberger, Phys. Rep.203 (1991) 345.

2. J.Bernabéu and P.Pascual, Electroweak Theory, Univ. Autónoma Barcelona, B-21399 (1981); S.M.Bilenky, Introduction to physics of electroweak interaction, Pergamon Press (1982).

3. J.Ashman et al., EMC experiment, Phys.Letters B206 (1988) 364; Nucl. Phys. B328 (1989)1.

4. E.N.Fortson and L.L. Lewis, Phys. Rep. 113 (1984) 289;
   U.Amaldi et al., Phys. Rev. D36 (1987) 1385;
   G.Costa et al., Nucl.Phys. B297 (1988) 244.

5. M.C.Noecker, B.P.Masterton and C.E.Wieman, Phys. Rev. Letters 61 (1988) 310.

6. S.A. Blundell, W.R.Johnson and J.Sapirstein, Phys. Rev. Letters 65 (1990) 1411.

7. F.Dydak, CERN-PPE / 91-14 (1991).

8. W.Marciano and J.L. Rosner, Phys. Rev. Letters 65 (1990) 2963.

9. G.Altarelli et al., CERN-TH 6028/91 (1991).

10. J. Bernabéu, Nucl.Phys. A518 (1990) 317.

11. Ya. B. Zeldovich, Zh. Eksp. Theor. Fiz 33 (1957) 1531 [JETP 6 (1957) 1184]; V.V.Flambaum and I.B.Khriplovich, Zh. Eksp. Theor. Fiz 79 (1980) 1656 [JETP 52 (1980) 835]; C.Bouchiat and C.A.Piketty, Z.Phys. C49 (1991) 91.

12. J. Bernabéu, T.E.O. Ericson and C.Jarlskog, Phys. Letters B50 (1974) 467; G.Feinberg and M.Y.Chen, Phys. Rev. D10 (1974) 190; A.N.Moskalev, JETP Letters 19 (1974) 216.

13. J.Missimer and L.M.Simons, PR-90-03, PSI (1990).

14. A.N.Moskalev and R.M.Ryndin, Sov. J.Nucl.Phys. 22 (1976) 71.

15. J.Bernabéu, J.Bordes and J.Vidal, Z.Phys. 41 (1989) 679.

16. P.Langacker, UPR-0451T (1990);
R.Casalbuoni et al., UGVA-DPT 1991/05-728.