Measurement of Parity Nonconservation in Atoms.

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1. – Introduction.

This lecture is an introduction to the subject of parity nonconservation in atoms. It is intended for the student or scientist who is not familiar with the field. The Colorado cesium experiment is described in detail, and we have attempted to present many of the technical details and considerations that led to the final experimental design.

The basic phenomenon discussed in this lecture is the parity-nonconserving (PNC) weak neutral-current interaction in an atom. As shown in fig. 1, this interaction arises from the exchange of a $Z_0$ boson between the electrons and the protons and neutrons in an atom. It can be contrasted with the much larger Coulomb interaction that arises from the exchange of a virtual photon between the electrons and the protons and predominantly determines the structure of the atom. Study of PNC in atoms began in earnest with the Weinberg-Salam-Glashow (WSG) electroweak theory which unified the weak and electromagnetic interactions and predicted a parity-violating neutral-current interaction.

In an atom, the parity-nonconserving Hamiltonian can be written as the sum of two parts [1]

\[ H_{\text{PNC}} = H_{1\text{PNC}} + H_{2\text{PNC}} , \]

\[ H_{1\text{PNC}} = \frac{G_F}{2\sqrt{2}} Q_w |\gamma_5| , \]

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where $G_F$ is the Fermi constant and $\gamma_5$ is the fifth Dirac gamma matrix acting on the electron. The largest part of the PNC Hamiltonian arises from the so-called weak charge, $Q_w$. This involves the axial-vector quark current and the contributions from all quarks simply add in similar manner to the weak electromagnetic charges add. The weak charge,

$$Q_w = 2(2Z + N)C_{1u} + (2N + Z)C_{1d},$$

involves the number of protons ($Z$) and neutrons ($N$), and the coupling constants $C_{1u}$ and $C_{1d}$[2]. These are two of the four fundamental coupling constants which characterize the neutral-current interaction between electrons and quarks. They are multiplied by the appropriate factors involving the number of up and down quarks, respectively. The total $Q_w$ is nearly proportional to the number of neutrons contained in the nucleus.

The $H_2$ PNC term, arising from the nuclear-spin-dependent or isospin contribution, is much smaller[3]. This contribution is about 1% of the total PNC, and, because the spins of the quarks in the nucleus tend to cancel, this nuclear-spin-dependent part is the smallest of which is the weak neutral-current contribution involving the quarks. The much larger part (≈ $\times 10$) is the nuclear anapole moment contribution[3]. This contribution comes from the more traditional weak interactions that take place in the nucleus which then couple to the electrons. We discuss this in more detail in sect. 7 below.

The primary interest in studying these interactions is to test the model of the elementary-particle interactions. This standard model is made up of the WSG electroweak theory plus quantum chromodynamics to account for the strong interactions. In the standard model (neglecting radiative corrections), the constants $C_{1u}$ and $C_{1d}$ are given by[2]

$$C_{1u} = \frac{1}{2} - \frac{4}{3} \sin^2 \theta_w, \quad C_{1d} = -\frac{1}{2} + \frac{2}{3} \sin^2 \theta_w.$$

The goal of atomic-parity-nonconservation work is to determine if these constants are accurately predicted by these formulae or if there are some corrections to these values which arise from «new» physics. By «new» physics we mean physics that is not included in the standard model[2].
To be more specific, we test the standard model and hence look for this new physics using the following prescription. 1) Obtain the value of $\sin^2 \theta_w$ from the mass of the $Z_0$ boson, which is now measured very precisely at CERN. 2) Use this value to calculate $Q_w$ as in eq. (3) plus the addition of small radiative corrections we neglected. 3) Compare the calculated value of $Q_w$ with what one obtains from atomic-parity nonconservation.

Before beginning a discussion of the very lengthy and difficult experiments and atomic theory which have been undertaken to determine $Q_w$, the obvious question is: «Why bother?». What is the motivation for doing this? The answer to this is that the standard model is almost certainly not the final answer. It has many undesirable features; in particular, there are a number of important quantities that have to be put in on a somewhat artificial ad hoc basis. Notable among these are the coupling constants, the masses of all the particles and certain basic symmetries such as the handedness of the neutrinos. None of these different quantities arise naturally out of the model, as one would expect from the ultimate theory. In addition, the standard model runs into trouble at high energies. As the energies become large compared to the mass of the $Z_0$ boson, the theory becomes divergent.

These features indicate that there must be new physics beyond the standard model. We note that many aspects of the standard model have only been tested to about 2%. This is not like quantum electrodynamics where the theory has been tested to nine decimal places and one is asking whether there might be some deviation in the tenth. Because the standard model has only been tested relatively crudely, there is still a reasonable likelihood that small improvements in precision may lead to fundamental insight into new physics. Finally, the standard model does not occupy an unchallenged position. Many extensions and alternatives have been proposed, and atomic PNC measurements are uniquely sensitive to the physics that many of these would produce. The reason for this unique sensitivity is that the two coupling constants, $C_{1d}$ and $C_{1u}$, are measured very poorly, or not at all, by the high-energy experiments. In addition, there are a number of radiative corrections to the standard model which have energy dependences. Since the atomic PNC is at low energy, the comparison of atomic-parity nonconservation with high-energy results is sensitive to such energy-dependent terms. All these features have combined to make atomic-parity nonconservation perhaps the only mainstream particle physics that is currently being done on a table top scale.

2. – PNC neutral currents in atoms.

Many of the early concepts of the effects of neutral currents in atoms were introduced by CURTIS-MICHEL in the 1960's[4]. However, activity in the field expanded rapidly after the papers of M. A. Bouchiat and C.
Bouchiat in 1974[1, 5] in which they considered this in the context of the WSG model.

The principal effect of this parity-violating interaction in an atom is to mix the $S$ and $P$ parity eigenstates so that the $S$ state is no longer a pure $S$ state but has a very small amount ($\delta_{\text{PNC}}$) of $P$ state mixed into it,

$$|S\rangle \rightarrow |S\rangle + \delta_{\text{PNC}} |P\rangle,$$

$$\delta_{\text{PNC}} = \frac{G_F}{2\sqrt{2}} Q_w \langle \gamma_5 \rangle_{\text{nuc}},$$

and

$$\delta_{\text{PNC}} = (5 \cdot 10^{-18})(ZC^1)(Z^2) \approx 10^{-11}.$$

This quantity $\delta_{\text{PNC}}$ involves the Fermi constant of the weak interaction $G_F$, the weak charge, $Q_w$, mentioned before, and an atomic matrix element $\gamma_5$ from the matrix element of the $\gamma_5$ evaluated over the nucleus. The evaluation is only over the nucleus because the $Z_0$ boson is a massive particle and the interaction is a short-range interaction. It is straightforward to estimate the approximate size of the mixing (eq. (7)). We have already indicated the weak charge is proportional to $Z$ multiplied by constants which are on the order of one. The matrix element, as pointed out by the Bouchiat, is proportional to $Z^2$. For a relatively heavy atom like cesium, the mixing then works out to be about $10^{-11}$, this very tiny scale that makes the experiment so difficult and fraught with the possibility of error. To put this in perspective, a ratio of one part in the same as the ratio of the diameter of a human hair to the diameter of the Earth.

In these experiments, the mixing of $S$ and $P$ states is observed as an electric-dipole transition amplitude between $S$ states with different principal quantum number $n$. Or, alternatively, between $P$ states with different $n$. In 1973, quantum mechanics one learns that such electric-dipole transition amplitudes are absolutely forbidden by the parity selection rule and, therefore, measures this electric-dipole transition amplitude, it is a measure of much $P$ state is mixed into the $S$ state.

The most direct approach to measuring this transition amplitude would simply take a laser, set its frequency in resonance with some $S \rightarrow S$ transition, and look for the excitation rate, trying to pick out the part due just to the electric-dipole contribution. In fact, in the days before Bouchiat's paper, this approach provided the best test of the conserving parity in atoms and set the limits at that time. However, the parity-conserving rate is proportional to the square of $\delta_{\text{PNC}},$

$$R_{\text{PNC}} = |A_{\text{PNC}}|^2 \propto \delta_{\text{PNC}}^2.$$  

Effectively, this means the oscillator strength for such a transition is $10^{-22}$ and, therefore, the transition rate is 22 orders of magnitude smaller.
normal allowed electromagnetic transition. With any conceivable experiment, this is an impossibly small transition rate, and will always be lost in the noise. Therefore, this approach is clearly unsuitable for achieving the level of sensitivity needed for investigating the weak neutral-current effects.

As discussed by Michel and the Bouchiats, a much more intelligent approach is to use a «heterodyne» or interference approach. In this case, the transition rate, $R_H$, is equal to the square of the sum of two transition amplitudes, one being the parity-nonconserving amplitude and the other being a much larger parity-conserving transition amplitude between the $S$ states:

$$R_H = |A_0 \pm A_{PNC}|^2 = A_0^2 \pm 2A_0 A_{PNC} + A_{PNC}^2.$$  

Here, $A_{PNC}$ is the quantity we are interested in, and $A_0$ is the parity-conserving amplitude. Notice that there is now a term $2A_0 A_{PNC}$, which is linear in the parity-nonconserving amplitude, and, therefore, can be large enough to measure. It is this interference term that we are interested in determining. This basic idea of mixing and detecting a small amplitude by having it mixed with a large amplitude is, of course, a very common one in physics and goes back to at least the early days of radio, if not before.

It is necessary to give some thought as to the phases of these amplitudes to make sure that such a parity-violating interference term can exist in any given experiment. One can work through all the mathematics and selection rules to find the appropriate conditions, but the basic point is that this is a parity-nonconserving term, and, therefore, one has to design an experiment which has a handedness built in. If the experiment has no such handedness, it will not be sensitive to a PNC effect. The second important question is, assuming the interference term exists, how does one isolate it from the much larger $A_0^2$ contribution to the rate? This is done by observing the modulation in the transition rate when the handedness is changed, or, in other words, a mirror reflection of the experiment is carried out. In that case there is a change $\Delta R_H$ in the transition rate due to the parity-violating $2A_0 A_{PNC}$ term reversing sign. Thus all the experiments we are going to describe determine a fractional change in the transition rate $\Delta R_H / R_H$. A key point is that the transition rate itself, $R_H$, is a very weak transition by normal standards, and the fractional changes we are interested in are very small. For example, in our experiment $R_H$ corresponds to an oscillator strength of $10^{-12}$ and the modulation $\Delta R_H$ is on the order of one part in $10^5$. Thus the basic scale of the experiment is set by the fact that one has a very weak transition, and, in order to see small modulations on that transition, there must be a very high signal-to-noise ratio. It is this issue which makes these experiments quite difficult, and is behind most of the design considerations that we will discuss. There is also a continuing tradeoff between a large $A_0$, which makes $\Delta R_H$ larger, but $\Delta R_H / R_H$ smaller, and a small $A_0$, which does the opposite.
3. – Experimental approaches.

The various experimental approaches that have been tried can be characterized according to what they use as an interference amplitude, $A_0$. The choice was to use an allowed magnetic-dipole ($M1$) amplitude for $A_0$. In this case, one drives a transition between two $P$ states with the same principal quantum number. This experimental approach has been used to measure parity violation in bismuth, lead and thallium. In this work, the actual observation is the optical rotation of the plane of linearly polarized light. This corresponds looking at a difference between the index of refraction for the left vs. right circularly polarized light. Obviously, such a difference reflects a handedness system. The basic set-up for such experiments is shown in fig. 2. The beam passes through a linear polarizer to ensure that its polarization is clean. It then passes through a vapor cell which contains the atom of interest and then finally goes into a second crossed polarizer. This second polarizer blocks out all the laser light unless its polarization has been rotated in the vapor, in which case some light passes through and can be seen at the detector. The actual experiments are somewhat more involved, but this is the basic idea. One then tunes the laser over the atomic transition and observes small rotation in the polarization as an increase in the light at the detector tuned on the transition.

Several steps have been taken to improve the signal-to-noise ratio and test for potential systematic errors, which have always dominated the uncertainty. First, to improve the signal-to-noise ratio, the polarizers are not slightly from perfectly orthogonal, and the incident polarization is modulated using a Faraday rotator. This latter step reduces the noise by shifting the detection bandwidth away from d.c. To eliminate sources of potential systematic errors, some or all of the following steps have been taken in the various experiments: 1) alternating between an oven containing atomic vapor and an oven with no vapor, 2) reversing the direction of the light through the vapor, and 3) careful fitting to the atomic line shape. The parity nonconservation is dispersion shaped and thus has quite a different dependence on laser frequency from the absorption.

The second approach is to use a Stark-induced $A_0$ amplitude for the interference. In this approach, one simply applies a d.c. electric field to the atom. The electric field mixes $S$ and $P$ states by the Stark effect, in a parity-conser-

![Fig. 2. Basic experimental set-up for optical-rotation PNC experiments.](attachment:image.png)
way as shown by

\[ |S\rangle \rightarrow |S\rangle + \varepsilon_E |P\rangle + \varepsilon_{\text{PNC}} |P\rangle. \]

One now observes an interference between the \( \varepsilon_E \) and \( \varepsilon_{\text{PNC}} \) mixing terms. This approach has been used to measure parity violation in cesium by th
groups at Paris and Boulder, and in thallium at Berkeley.

Each of these approaches (allowed M1 and Stark induced) has certain advan
tages and disadvantages. While these often get quite technical, we can summa
rize some of the more notable features. First, let us consider the M1 approach.
It can be characterized by the fact \( A_0 \) is quite large, which makes \( \Delta R/R \) small
However, since the transition rate itself is large, the statistical signal-to-noise
ratio can be quite good. Another advantage of these experiments is that they
are relatively simple. The drawback to this approach is that, because the frac
tional modulation is quite small (~ 1 part in \( 10^8 \)) and there are relatively few re
versals to isolate the PNC component, the systematic errors are major
problems.

This can be contrasted with the Stark interference approach. Here \( A_0 \) is
small and the problems generally are statistical signal-to-noise ratios. There are
many reversals to isolate the parity-nonconserving effects and suppress sys			
tematic errors and the fractional modulation is relatively large. Therefore, sys					
tematic errors are much less of an issue. Because these experiments involve
extra applied fields and more reversals, they tend to be rather more compli
cated, however.

The other important experimental issue is the choice of atom. Here there are
a number of considerations that come into play. First, one wants a heavy atom
because of the \( Z^3 \) dependence (eq. (7)) of the mixing. Second, one needs an
atomic transition between states with the appropriate quantum numbers. If one
is using the magnetic-dipole interference, this needs to be a \( np \)-to-\( n'p' \) transi
tion. If one is using the Stark-induced approach, either an \( nS \rightarrow n'S \) or an \( np \rightarrow
\rightarrow n'p \) will work. In this case it is desirable to have different initial and final \( n \)'s to
suppress the magnetic-dipole amplitude, which can introduce systematic errors
and background noise. Also, in all cases it is desirable to have the transition at a
wavelength where a good tunable laser is available to excite the atom. If one is
using fluorescence detection (as in Stark-induced experiments), it is important
that the fluorescing light be at a wavelength that is significantly different from
the excitation light. Otherwise, the scattered excitation light will give an over					
whelming amount of background. Another consideration is that the atoms must
remain isolated rather than forming molecules, where the PNC effects are ob					
sured. The final important issue, and at this stage of the field clearly the most
important, is the accuracy with which the \( \gamma_5 \) matrix element for the atom can be
calculated. This is important because one is now worried about precision tests
of the standard model rather than simply detecting parity nonconservation. We

will discuss this issue in more detail later. Here we will simply summarize for alkali atoms such as cesium, francium and rubidium, these calculations be done with an accuracy on the order of 1%. In thallium calculations the accuracy is on the order of 5%, while for systems such as bismuth and lead, many valence electrons, the accuracy of calculations of the $\gamma_5$ matrix element relatively low.

Before discussing the experiments in more detail, we might mention other possibilities for PNC experiments which have been proposed but not tried out. Among other interference amplitudes that can be used, there are very few that have been seriously discussed. There have been proposals by Anderson [6] to interfere a two-photon-allowed transition amplitude with one-photon PNC amplitude. In this case, the interference is sensitive to relative optical phases. This can be both an advantage and a disadvantage, ever, and this approach has not been seriously pursued up to this time.

A number of other experiments have been proposed, and a few attempts involving other atomic species. Probably the most notable is the use of hydrogen. In this experiment, the idea was to observe the mixing of the 2S and 2P states of hydrogen. Although $Z$ is 1, and therefore, the PNC matrix element is quite small, this is largely offset by the mixing energy denominator which is nearly zero for these two states. Hence, the actual $\delta_{\text{PNC}}$ for $n = 1$ hydrogen is nearly the same as that for heavy atoms. Because of this there were number of experimental programs initiated to study PNC in hydrogen. However, the great problem with the hydrogen case is systematic errors. Electric fields, which can also cause mixing of the 2S and 2P states, are a field by the same near-zero energy denominator. As a result, the systematic errors, relative to the PNC signal, are enhanced compared to heavy atom by factor. Effectively, this means that, instead of needing to worry about mV stray fields, one has to worry about nV/cm fields. This is a nearly impossible problem and, therefore, to our knowledge, all the experiments on hydrogen PNC have now been abandoned.

Another set of experiments which have been proposed and pursued to degree involve the use of ions or muonic atoms for studying PNC. In these es, the overlap of the electrons at the nucleus is much larger than for a neutral atom, and, therefore, $\delta_{\text{PNC}}$ can be relatively large. However, this is more offset by the fact that the sample size is very small. At the present time technology is not available to produce large enough samples to allow meaningful PNC measurements. However, this is very much a function of technology which is likely to change.

We would now like to discuss the first generation of Stark interference experiments which were carried out in Paris [7] and Berkeley [8]. The schematic for these experiments is shown in fig. 3. A circularly polarized beam is sent into an atomic-vapor cell and excites the transitions of interest in the presence of a d.c. electric field. The transition rate is monitored by ot
ing the fluorescence as the atom decays back to the ground state. Since the Zeeman transitions are not resolved, the parity-nonconserving interference term cannot be observed directly in the total atomic-transition rate, but it does cause a polarization of the excited state. This polarization is detected by looking at the degree of circular polarization of the fluorescence light. These experiments were successful at detecting a small circular polarization and hence a parity violation. It was measured with a fractional uncertainty at a level of (10 ± 20)%.

Before beginning a lengthy discussion of the design considerations of the Colorado experiments, we would like to briefly review the lessons we have learned from these pioneering first-generation optical-rotation and Stark interference experiments. The two main lessons from the optical-rotation experiments are that the systematic errors are very serious and must be considered in great detail, and the atomic-structure issue is crucial if one is interested in a precision test of the standard model. These considerations made us decide to go with the Stark interference approach and use cesium, since Cs is a heavy alkali atom. However, there were also several important lessons from the Stark interference experiments. The first was that the signal-to-noise ratio was disappointingly low, predominantly due to problems from background signals. These signals can arise from sources such as blackbody radiation, molecules, collision-induced transitions and scattered light. Also, in these cell experiments there were a very large number of nonresonant atoms, atoms that were not being excited by the laser beam and did not contribute to the PNC signal. These atoms could give rise to background noise or systematic errors, which is clearly undesirable.

Finally, a nonobvious design feature is the issue of experimental flexibility. In such cell experiments, it is very difficult to build a new cell and, therefore, one is quite limited in how easily one can change the experiment in response to new data or ideas.
4. – Design concepts of Colorado experiments.

After reviewing the lessons from the work described in the previous we had three major concerns. The first was to have flexibility in the ap so that the experiment could be easily modified. The second concern wa crease the statistical signal-to-noise ratio (S/N), and the third was to potential systematic errors.

To improve the S/N, we set out to achieve high detection efficient minimal background. One way to achieve this is to have a modulation dir the transition rate rather than the polarization of the excited state. An magnetic field would allow one to do this by resolving the different $m$ le is particularly useful in conjunction with a collimated atomic beam I much weaker magnetic fields are required.

There are also several other desirable features to an atomic beam. I has rather few nonresonant atoms and molecules and allows high detecti ciency. Second, the beam can be turned off and the system opened changed quickly. The major drawback is that the actual number of atoms one has is significantly lower than in a cell. However, we felt t other factors would more than offset this, which has proved to be th

The primary feature for controlling potential systematic errors is to an experiment with many mirror reversals. However, in addition, it is tant to have ways to quickly measure potential systematic errors will precision.

It is quite easy to understand how the addition of a magnetic field in these experiments. The $A_{\text{PNC}}$ amplitude is proportional to the magnetic number $m$, while the Stark interference amplitude $A_E$ is proportional the absolute value, $|m|$. Therefore, the sum over all $m$ of the product $A equals zero. This arises because of the requirement of time-reversal co tion. However, if one only looks at transitions between individual $m$ then the product $A_{\text{PNC}} A_E$ is not equal to zero, and will contribute to the tion rate. There are two ways one can experimentally observe transiti between individual $m$ levels. The first way is to simply apply a large enou netic field and use the Zeeman effect to resolve the different $m$ levels. Ti laser will only excite transitions from a single $m$. This approach was used Colorado 1985 and 1988 experiments[9, 10], and in the Drell and Comm lium experiment[11] at Berkeley. The second option is to simply optically all the atoms into a single $m$ level. This has the advantage of utilizing the much more efficiently, but it does make the experiment more complicate approach is being used in the current experiment at Colorado.

The relevant energy levels of the cesium atom are shown in fig. 4. ' ground state has two hyperfine states with total angular-momentum qi numbers $F = 3$ and $F = 4$. The $7S$ excited state also has $F = 3$ and $F = 4$ fine states. In the presence of a magnetic field, each of these levels the
up into the different $m$ levels, and one excites $6S \rightarrow 7S$ transitions between these independent $m$ levels with 540 nm laser light. The excitation of the 7S state is monitored by observing the fluorescence produced primarily at 852 and 894 nm when the 7S state decays by an allowed transition to the 6P and then to the 6S states. In fig. 5, we show the theoretical spectrum of the $6S, F = 4 \rightarrow$

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**Fig. 4.** - Cesium energy level diagram.

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**Fig. 5.** - Theoretical spectrum of the $6S, F = 3 \rightarrow 7S, F' = 4$ transition of cesium in a weak magnetic field. **a)** The pure Stark-induced rate ($A_E^3$), **b)** the $A_E A_{PNC}$ interference terms multiplied by $10^6$. 

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→ 7S, F′ = 3 transition in a weak magnetic field. Figure 5a) represents the Stark-induced rate that is proportional to A2; figure 5b) represents the A interference terms. This illustrates how the parity-nonconserving contribution adds to the rate for the +m levels, while subtracting for the −m levels.

A general layout of the experiment is shown in fig. 6. A collimated beam is intersected by a laser beam which excites the 6S → 7S transition in the interaction region, there are three perpendicular vectors defining a coordinate system. These are a d.c. electric field, E, a d.c. magnetic field, B, a angular momentum σ of the laser photons. The 6S → 7S excitation rate is measured by observing the subsequent fluorescence. In fig. 7, we show a spectrum observed when we sweep the laser over one of the hyperfine transitions. This shows eight lines of the Zeeman multiplet corresponding to excitation of different m levels. It is quite obvious that there is a rather peculiar line shape. The explanation of this line shape is somewhat complicated. It involves the off-resonance a.c. Stark shift in combination with the small D shifts. (This has been analyzed and explained in ref.[12].) The values of power and cesium beam divergence for this figure make the line shape less peculiar. For larger or smaller power, or less beam divergence, the line becomes smoother and more systematic. In any case, this strange line shape is a minor anomaly and has no real effect on the parity nonconservation measurements.

To summarize the basic experiment, we set the laser frequency on one peak of the Zeeman multiplet and look for a change, typically a part in the transition rate when the parity or handedness of the experiment is reversed. There are many different ways to reverse the handedness; most of them can be seen simply by considering the coordinate system defined by E, B. Anything that reverses the handedness of this coordinate system will reverse the sign of the parity-nonconserving term. Equivalently, this can be described as making a mirror reversal with the mirror oriented in three different ways. Thus the three reversals are E → −E, B → −B, σ → −σ (right circular

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Fig. 6. General layout of the Boulder cesium PNC experiment. The coordinate system is defined by a d.c. electric field, E, a d.c. magnetic field, B, and angular momentum, σ, of the laser photons.
Fig. 7. – Experimental spectrum observed when the laser is scanned over the 6S, \( F = 4 \rightarrow 7S, F' = 4 \) transition of cesium in a 70 G magnetic field.

ized light to left circular polarized light), and, finally, \( m \to -m \). This \( m \) reversal is carried out by simply changing the frequency of the laser so that it moves from one Zeeman peak to the symmetric Zeeman peak on the other side of the multiplet.

The presence of four independent reversals is crucial for suppressing systematic errors. Having many reversals provides a redundancy which tests the signal observed truly is parity violating. It greatly suppresses potential systematic errors which might mimic PNC under one or two reversals, but is unlikely to pass the test of all four.

While the experiment is rather simple in principle, in fact a tremendous amount of time and effort has gone into optimizing the apparatus to achieve the necessary signal-to-noise ratio. The next section will discuss in detail these signal-to-noise issues, and how they have led to the construction of the apparatus in its present form.

5. – Details of apparatus.

5.1. Signal and noise analysis. – The majority of the time spent on this experiment has been used to achieve the necessary signal-to-noise ratio. The total detector current is given by

\[
R \propto \{(\text{number of atoms})(\text{laser power})(\text{detection efficiency})\} \cdot \left[ E^2 \pm 2E \delta_{\text{PNC}} \right] + \text{background}
\]

\[
S = \Delta R \propto \left[ 4E \delta_{\text{PNC}} \right] \text{ and is } 10^{-5} \text{ to } 10^{-6} \text{ of } R \text{ in this experiment.}
\]
The main signal depends on the number of atoms, the laser power a
tection efficiency multiplied by the term in brackets involving the ele
Of course, the largest electric-field term is the $E^2$ pure Stark compo
there is the much smaller PNC interference term which is linearly pi
to the electric field. The actual signal of interest, $S$, is the change in
tor current when the fields are reversed. This is proportional to the
atoms $\times$ laser power $\times$ detection efficiency $\times 4E^2_{\text{PNC}}$. The noise, h
made up of several parts which we will characterize according to th
dence on the electric field $E$ as given by

\begin{equation}
N = N_{\text{bk}} + \text{(detector noise + fluctuations in back)} + N_{\text{sh}} + \text{(shot noise on } R, \alpha) + N_{\text{tech}} \text{ (technical noise on } R, \alpha) \end{equation}

First there is the background noise, which is independent of $E$. This is
fluctuations in the signal due to the fundamental detector noise or
grounds due to scattered light, room lights, etc. The second term is the
shot noise fluctuations on the total detector current. Since this is propor
tional to the square root of the current, it is linearly proportional
ally, the third term, what we call the technical noise, is propor
times some fluctuation fraction $f$. Here, $f$ can be due to changes in a lier of things, for example, laser power, laser frequency and atomic-be
sity. Because of this dependence of the noise on $E$, one has an inter
dependence of the signal-to-noise ratio of $S$ on the electric field. This is
fig. 8.

For an ideal case, where ideal is rather low technical and backgro
one would have the signal-to-noise ratio rising linearly with electric fi
region where one is dominated simply by background noise. At high
shot noise, which is proportional to $E$, begins to dominate. At this
signal-to-noise ratio is independent of $E$ since both the signal and the

![Fig. 8. Dependence of the signal-to-noise ratio on the electric field.](image-url)
proportional to $E$. Finally, at even larger $E$ fields, the technical noise with $E^2$-dependence will dominate. When this occurs, the signal-to-noise ratio falls as $1/E$. In this ideal case one could then set the electric field anywhere in the region where the shot noise dominated and achieve the same signal-to-noise ratio. In fact, this is not the case in the real world.

In these experiments, we have always had a situation more like that shown in the lower curve of fig. 8 where the background and technical noise are large. Here, the signal-to-noise ratio rises with the electric field until it reaches some peak, above which the technical noise dominates. There is no actual flat region where the shot noise dominates, and, therefore, there is always some optimum electric field. We work hard to push this peak up until it is fairly close to the shot noise limit, and then there is little to be gained by further improvement. At this field, typically the background noise and the technical noise are equal. The signal-to-noise ratio works out to be roughly proportional to the square root of the number of atoms times the laser power times the detection efficiency. In order to reach the shot noise signal-to-noise ratio, however, it is necessary to make the background noise small and to minimize the technical noise by making the quantity $f$ small. The former means keeping the detector noise and scattered-light noise low, while the latter requires a number of parameters be highly stabilized. The primary feature that makes these experiments so difficult is the fact that achieving this to the level necessary to make high-precision PN measurements requires pushing many different technologies to the state of the art or beyond the state of the art. This is different from many experiments where most of the experiment is based on well-developed technology and there are only one or at most two different technologies that one has to work on very hard.

Since the signal is proportional to the number of illuminated resonant atom the laser power and the detection efficiency, it is clearly important to make at least three of these large. We will discuss efforts to maximize each of these factors separately.

5.2. Atomic beam. – In this experiment for an atom in the beam to be «useful» it has to meet slightly different requirements from most atomic beam where the only requirement is either flux or density. In this case a long thin laser beam is exciting a narrow (2.8 MHz) transition. Thus atoms which have divergence such that their velocity parallel to the laser beam gives a Doppler shift larger than 2.8 MHz are out of resonance with the laser and will not contribute to the signal. The divergence in the perpendicular direction affects the density but does not cause a Doppler shift. Thus the three basic consideration are: the density of atoms in the beam, the length of atomic beam intercepted by the laser and the velocity of the atoms parallel to the laser beam (the divergence in one direction).

If one starts with an atomic vapor and wants to make a collimated beam ou
of it, the simplest approach is to let it effuse out of a tube of length \( L \) or \( R \). This tube is the output nozzle on an oven. In this case the intensity in the forward direction will increase as the pressure of the vapor is increased, the mean free path of the atoms in the vapor becomes less than \( L \). At this point the atoms coming out of the oven simply form a cloud, the divergence of the beam increases, and the number of atoms in the acceptable divergence angle will stay constant or decrease. To improve on this, we have used a technique that had been successful for helium, but which had never been successfully applied to cesium. This technology is a glass capillary array nozzle. Such an array is made up of thousands of very tiny tubes, typically 10 \( \mu \text{m} \) in diameter and 1 cm long. In order to increase the width of the beam, we use a section of the array which is 0.4 cm by 2 cm so that it provides a collimated ribbon of cesium. This provides a more intense beam with the appropriate divergence angle than can be provided by a single tube or slit.

We have found that operating a glass capillary array at the pressure of the mean free path is equal to \( L \) does not give the most intense possible beam. We have improved on this by using a glass capillary array as a nozzle in an oven which is operated at a vapor pressure which substantially degrades the intensity of the beam. About 2 cm from the array we place a multislit collimator made of a large number of thin metal veins which provides collimation only in the laser beam direction (fig. 9). In between these two collimators, we have placed cryogenic panels which pump away the background Cs vapor. To avoid the Cs\(_2\) dimers, which cause a very noisy background signal, we keep the glass capillary array about 100 °C hotter than the rest of the oven. We have been able to produce a state-of-the-art atomic beam in this manner, the beam intensity was lower than what we hoped for. The intensities we achieved were based on results that had been achieved with helium, but we achieved less than 1/10 of this goal.

53. Laser power build-up. — Fortunately, we have been able to compensate for lower than anticipated beam intensity by having more laser power available than we had originally anticipated. The 540 nm laser light which excites the...
comes from a c.w. dye laser. The transition rate is very small, so only a small fraction ($< 10^{-11}$) of the light is absorbed as it passes through the atomic beam. This means that one can gain tremendously by reusing the light. There are two common ways to reuse light when looking at weak transitions. The first technique, which has been by far the mostly widely used, is to have a multipass system as shown in fig. 10. Here the light comes in and is reflected back and forth between two mirrors and then finally leaves without ever overlapping itself. This technique was used in the Bouchiat cesium experiment, and we initially tried to use it for our experiment, but we discovered it had some disadvantage. First, the number of bounces one can get is limited to the size of the mirrors divided by the size of the laser beam, since the beam occupies a different spot on each bounce. With a reasonably sized mirror, one has a limit of a few hundred bounces. Also, the interaction region becomes quite extended, which is a very serious limitation if one wants efficient detection. This is because imaging fluorescence from an extended region is always much less efficient than if the fluorescence is very localized. Finally, there is also a somewhat subtle effect involving systemic errors due to birefringence of the mirrors which is difficult to deal with in this configuration because one can only observe the average effect of many reflections.

As a result of these difficulties, we switched to a different configuration shown in fig. 11. This is known as a «build-up cavity» because a resonant Fabry-Perot interferometer builds up the laser power between the two mirrors. The most important issues in a build-up cavity are the losses and transmissions of the two mirrors. To optimize the build-up, the transmission of the output mirror should be as low as possible. The transmission of the input mirror should be between one and two times the round-trip loss in the cavity. Of course, any realistic book will provide the equations you need to calculate the exact optimum. Physically, what is happening is that, if the input transmission is much lower than the losses, the light will die out inside the cavity faster than it builds up. On the other hand, if the transmission is much larger than the losses, then the

$$P_0 \rightarrow \frac{P_{in} - P_0/(1-T_1)}{T_1 - 2 \times \text{loss}}$$

$$T_1 - 2 \times \text{loss} \quad T_2 = 0$$

Fig. 11. – Laser power build-up cavity. $T_1$ is the transmission of the input mirror, $T_2$ is the transmission of the second mirror.
light will leak back out the way it came in before it can reflect back as many times. From this argument it is easy to see why the optimum transmission should be comparable or slightly more than the total losses due to reflection and scattering. In this case, one achieves a traveling-wave power in cavity which is equal to the incident power times the build-up factor \( B = T_1^{-1} \), where \( T_1 \) is the input transmission. This makes it clear that the one obtains is limited entirely by the mirror losses. Over the course work, as we have improved mirror coatings, we have progressed by a build-up factor of 100 to 1000 and, in our latest work, to 15,000. With 0 incident power, we have 4.5 kW of circulating power inside the cavity.

Unfortunately, this tremendous enhancement in the laser power, in the signal, does not come without a price: laser stabilization. This is a major disadvantage of the build-up cavity relative to a multipass cavity where there is no resonance. For a build-up cavity, the laser must be stabilized to cavity resonance, which is only about 8 kHz wide for a build-up factor of 15,000. However, it is not sufficient to simply stabilize the laser to 8 kHz. If the cavity width varied by 8 kHz, our \( 6S \rightarrow 7S \) signal would be varying by nearly 100% the desired measurement accuracy is about 1 part in \( 10^7 \). Therefore, it is necessary to have the laser stabilized to a small fraction of this 8 kHz cavity resonance. Since the typical free-running dye laser has short-term frequency instability of about 10 MHz, this puts severe demands on laser stabilization. In general, even if the laser is locked to the cavity, should the cavity move relative to the atoms’ transition frequency, there will also be large changes in rate. Hence, the cavity must be stabilized to the atomic resonance to within a small fraction of the atom’s natural linewidth. This requires considerable stabilization to the mechanical and thermal stability of the cavity. Ultimately, stabilization is done with servo loops; we will discuss some of the control servo systems used in this experiment in more detail in subsect. 5.6.

5.4. Detection. — The third important factor that determines the sign is the detection efficiency of the \( 6P \rightarrow 6S \) fluorescence which is produced as a result of the \( 7S \) excitation. As mentioned earlier, the more localized the point source, we do nearly as well by having a line source of fluorescence, is defined by the region where the thin laser beam intersects the ribbon beam (fig. 12). We use a cylindrical mirror to focus the light from this interaction region onto a long, narrow detector. Essentially we have a two-dimensional system, which allows us to obtain a rather high collection solid angle of \( 2\pi \) sr.

A major difficulty is that the fluorescence photons must be collected in the same region where we must apply very carefully controlled electric fields. The only practical way to accomplish this is with parallel field plates, and the photons must pass through the electric-field plates. This required th
Fig. 12. – Fluorescence collection and detection.

Development of appropriate transparent, highly conducting coatings for glass. The development and stability of these coatings has been a major headache this experiment.

Silicon photodiodes operating at liquid-nitrogen temperature are used for detection. Although it is generally believed that a photomultiplier is better for very sensitive detection, this is not true for relatively large signals. Silicon photodiodes have nearly 100% quantum efficiency at 850 nm, but a photomultiplier detection efficiency is only about 10% or less. Of course, a PIN photodiode has no gain, whereas a photomultiplier has gain which introduces very little noise. However, if the signal is large enough that the photon shot noise is comparable to the detector noise, then the photodiode is superior due to its quantum efficiency. Thus it is quite important to have the photodiode noise small. One way to accomplish this is to have a small diode, but this is incompatible with having a large detection solid angle. However, even with a large (1 cm) photodiode, we have been able to achieve a detector noise comparable to about $10^4$ photons/s$^{1/2}$ by cooling it to liquid-nitrogen temperature. It is also necessary to use a very low noise op-amp with a large feedback resistor in the current-to-voltage amplifier so that amplifier noise is not a significant limitation. This means that the photodiode actually gives the best signal-to-noise ratio for signals larger than about $10^7$ photons/s. Clearly, if one is interested in detecting effects of a part in $10^6$, it is necessary to have far more than $10^7$ photons/s.

Having efficient detection with low detector noise is only half the detection problem. The other half concerns the noise arising from background signal. The major source of background is the scattered laser light reaching the detector. This is a large problem because there are $10^{22}$ green laser photons in the power build-up cavity for every $10^8$ infrared photons produced. Thus an isolation factor of $10^{13}$ is required! This isolation is achieved partly by geometry simply putting in light baffles which absorb any green scattered light and shield the detector. Of course, this is never perfect, so we also use colored glass filters to block out the green light but pass the infrared photons. Colored glass filters are not as selective as interference filters, but they have a large acceptance angle, which allows much better detection efficiency. When we first looked at filter specifications, we thought there would be no problem in obtaining a filter that suppressed green light and transmitted near i.r. light. How
ever, this was not the case. The specifications do not mention that green tons cause i.r. fluorescence in the filter at a part in $10^5$ or $10^6$, which is transmitted through the filter. We spent considerable time studying this process and investigating many different filters. We found the conversion efficiency varies widely, but eventually obtained filters that worked quite well. It is interesting to note that the best filter was not one of the standard colored filters, but in fact was the red plastic in the back cover of the Schott book.

5.5. Technical noise. – These various signal and background issues led to the final design of the 1985 and 1988 experiments shown in fig. 13. Light from a dye laser passes through optics, which control polarization very precisely, then enters the power build-up cavity that is inside a vacuum chamber. Necessary to have the cavity mirrors inside the vacuum chamber because there are no windows with low enough birefringence and loss to be put inside the chamber. The dye laser is frequency stabilized to the PBC with an electronic loop and its intensity is stabilized with another servo loop. The light inside the power build-up cavity intersects the atomic beam in a region of d.c. electromagnetic fields. The details of the interaction region are shown in fig. 13.

Once this apparatus is assembled, the prime concern is reducing the technical noise and eventually reducing the size of the systematic errors. The technical noise must be reduced to a level where the modulation fraction, $f$, is less than or equal to $10^{-5}/\sqrt{\text{Hz}}$ ($N_{\text{tech}} = f \times R$, where $R$ is the signal size). There are many processes that can produce technical noise at this level, including several that would not be encountered in a more typical experiment. For example, we have seen fluctuations in the atomic beam at this level. We found that the standard atomic beam is actually very quiet, but fluctuations of the pressure in the vacuum chamber of $\gtrsim 10^{-8}$ Torr cause noise on our signal. We found that such pressure fluctuations come from the release of small quantities of gas.

![Diagram](image-url)

Fig. 13. – Overall schematic of the apparatus for the 1985 and 1988 Boulder cesium clock experiment.
bles of gas. These can be eliminated with proper design of the vacuum pumping system.

A more obvious source of technical noise is the laser power fluctuations inside the power build-up cavity. These can come from fluctuations in the laser power itself or fluctuations in the laser frequency. The latter is the more difficult to correct; as we mentioned earlier, an 8 kHz change in the laser frequency relative to the cavity resonant frequency, leads to \( f = 1 \). Therefore, considerable effort must be made to stabilize both the frequency and the intensity of the laser. This is done with several servo loops. First, the laser is locked to a build-up cavity using the Pound-Drever frequency modulation technique. This lock has good signal-to-noise ratio and a rather fast response time, so that it can correct the errors in a time as short as 1 \( \mu \)s. However, the length and hence the resonant frequency of the build-up cavity can fluctuate because of inherent instabilities in the mechanical design and vibrations of the environment. To avoid this we lock the build-up cavity to a very stable reference cavity. Because this cavity does not have an experiment inside it, it can be much more rigid. The cavity-to-cavity lock also has good signal-to-noise ratio, but has a relatively slow response (400 Hz). This response is limited by the fact that we cannot move the build-up cavity mirrors very fast without introducing strains that cause birefringence and lead to systematic errors. Although the reference cavity has good short-term stability, it is subject to slow drifts and, therefore, must be locked to the atomic-transition itself. This lock has relatively poor si
nal-to-noise ratio because it relies on the atomic-transition rate, which is this experiment. Fortunately, the drift rate of the reference cavity is so one does not require the high signal-to-noise ratio necessary for a very fi: vo response. The signal-to-noise ratio limitation on this lock is the real- use the reference cavity rather than locking the frequency of the build-i- ity directly to the atoms.

5'6. Servo systems. – Servo-control systems are very important to t: many other «frontier» experiments of the sort discussed elsewhere in the: uma. Because they are often not part of a general physics education, v: provide a brief introduction to servo theory, which will allow one to gra: key issues involved in the stabilization of systems. A basic servo loop is: in fig. 15. The laser has some noise, for example, the table vibrations: move the mirrors and thus cause the output frequency to vary. One can: pare this output frequency with some reference, such as a stable optical: and from this obtain an error signal that indicates how far the frequen- moved away from what is desired. The basic error signal is then modifi: electronics which provide gain and some type of compensation and the re: a feedback signal that is sent back to the laser. This feedback is negative: cancels the errors introduced by the table vibration and shifts the fre: back to where it is supposed to be. With this feedback, the error signal, E: comes \( E_f = E / \text{gain} \) in the limit of gain \( \gg 1 \), where \( E \) is the error signal with feedback. Thus one wants to make this gain large in order to push the error signal as close to zero as possible, which is equivalent to forcing the laser frequency to be the same as the reference frequency.

The key point in designing any servo loop is that there are always tin: lays in the system, and much of the design is based on dealing with these delays. This can be best understood by considering the laser mirror. The m: is attached to a piezoelectric transducer (PZT) which will stretch it when: age is applied. Although the PZT is a hard ceramic and the rest of the m: solid metal piece, on the scale of frequency errors and corresponding dist:

Fig. 15. – Basic servo loop.
relevant here (atomic diameters or less), there is no such thing as a stiff mou
In fact, at these atomic scales the stiffest PZT is incredibly spongy and can cor
press many wavelengths. Therefore, a PZT can be best visualized as a sprir
with the mass of a mirror mounted on it, as shown in fig. 16. We then move t
back end of the piezo to try to correct the position of the mirror and compensa
for random vibrations. If one considers how a mass that is attached to a wall t
a spring responds to sinusoidal motions of the wall as a function of frequenc
this is just a driven harmonic oscillator. For frequencies below the resonant fr
quency, the response is in phase with the drive. However, above the resonant f
requency of the mass-spring system, there is a 180° phase lag and the mirror
moves in the opposite direction of the driving force. This shows the primar
problem encountered in designing any servo system; if the amount of feedback
is the same amplitude and phase relative to the error signal for all frequencie
the servo works fine for correcting for errors that fluctuate at frequencies lowe
than the resonant frequency. However, above the resonance frequency, the
180° phase shift causes positive feedback and the system becomes an oscillator
the feedback gain is greater than 1.

This is obviously unacceptable. A straightforward solution is to make the
gain smaller than 1 for frequencies above the resonant frequency and make
larger than 1 for lower frequencies. This gain is produced simply by having th
compensation electronics include a simple low-pass filter. The normal way t
operate a servo system containing such a filter is to turn up the gain until it jus
starts to oscillate at the 180° phase shift point, and then reduce the gain slightl
so that it stops oscillating. This provides a stable servo system. Notice that th
gain at low frequencies is set by the response (phase lag) of the system at high
frequencies. This is characteristic of any servo system.

Here we have presented the simplest possible compensation. In a more ad
vanced servo design, one would put in more elaborate compensation involving
electronic circuits that change the phase shifts and gain with frequency. With
such systems one can optimize the gain at particular frequencies where the
noise might be especially large or where one wants to have the system be par
particularly stable, such as at the frequency where the data are being act
other reasons for more elaborate compensation are to improve the transi
response so the system can recover more rapidly from a sudden shock. In t
main of this section we provide a few more examples of relatively stu
forward compensation and how one can deal with various kinds of phase l
systems.

Several methods have been developed for designing servo systems
quency response methods, the root locus method and state space method
method that one chooses depends on the requirements of the servo design
as transient response and steady-state error. We will describe frequen
response methods from an experimentalist's point of view. We have chose
method because it is very easy to measure a system's frequency respon
modern signal analyzers.

The basis of the frequency response method is the Bode plot. The Bod
shows a system's gain and phase as a function of frequency. Both the syst
be controlled and the compensation have characteristic Bode plots. When
designing a servo, one first measures the frequency response of the system
controlled and then designs a compensation circuit that tailors the open
frequency response to provide the desired control. It is important to ke
mind that high gain at frequencies where the phase is less than $180^\circ$ pre
good control, but low gain is required at frequencies where the phase is ge
than or equal to $180^\circ$.

As an example of how to tailor the frequency response of a system with
ensation, consider a simple harmonic oscillator. Many physical systems c
modeled as damped harmonic oscillators. The Bode plot of a harmonic osc
is shown in fig. 17. There are two different goals one can have in trying t

![Bode plot for damped harmonic oscillator.](image)
trol this system: i) to minimize the steady-state error, i.e. to have a large gain; or ii) to maximize the bandwidth of the servo, i.e. to provide damping the resonant peak with feedback.

As we mentioned previously, the simplest way to prevent oscillation is make the gain smaller than 1 at frequencies at which the phase shift is greater than 180°; a Bode plot of such compensation (an integrator, or a low-pass filter is shown in Fig. 18a). Figure 18b) shows the resultant frequency response of the oscillator-compensation system. Note that the gain at low frequencies has increased, but the phase shift has reached 180° at a lower frequency. We no longer have a larger d.c. gain but a smaller bandwidth. The controlled oscillator w.
lock to the reference signal well at d.c., but will have a slow transient response and more noise at higher frequencies.

Suppose instead one compensates by adding a «phase lead» (e.g., a differentiator) to the compensation so that the phase shift of the oscillator-compensation system has not yet reached 180° at the resonance. This allows the resonance to be artificially damped. Figure 19 shows the resultant frequency response harmonic oscillator compensated with a phase lead. Note that, comparing Figure 18, the gain near the resonant frequency is large, and the d.c. gain has increased. This system will have a faster transient response, but more error locking to the reference signal at d.c.

There is in general a trade-off between bandwidth (fast transient response of a servo and d.c. gain (small steady-state error) of a servo. One way to think about this is in terms of integrators and differentiators. An integrator will generally provide less steady-state error because it has «memory» to make accurate adjustments at low frequencies. On the other hand, fluctuations that are fast-pared to the integration time will be «washed out», and, as a result, the bandwidth of the system will be reduced. Differentiators predict the future performance of the system by looking at the slope of the error signal, and, therefore, increase the bandwidth. However, because a differentiator is compensating for future fluctuations, it can slightly over or under compensate, leading to steady-state accuracy.

Systems with more complicated frequency responses than that of a single harmonic oscillator can be controlled by extending these ideas. A compensating circuit's phase and gain characteristics are tailored to provide the best compromise of bandwidth vs. d.c. response.

We now return to discuss how this stabilization is applied to the actual system.

Fig. 19. – Bode plot for harmonic oscillator plus phase lead compensation.
experiment. The power build-up cavity is stabilized by moving the mirrors vs piezoelectric transducers. Then laser frequency is locked to the power build cavity by a combination of elements, most of which are standard in c.w. lasers. First, we have a rotating plate on a galvanometer that changes the optical length of the laser cavity. This has a rather slow response, but a large namic range. Second, one of the mirrors is mounted on a PZT. This can change the cavity length with a frequency response extending to about 50 kHz. The fastest feedback is provided by an electro-optic modulator. This has a gain frequency of about 2 MHz, but can only correct for rather small errors in the frequency. In addition to the frequency stabilization, to stabilize the output power inside the power build-up cavity, we sense the light transmitted by the output mirror and hold it constant using acousto-optic or electro-optic modulators to control the incident laser power. One significant difficulty in this experiment is the fact that the power transmitted by this mirror does not seem to exactly proportional to the power inside the cavity at the parts in 10^6 level. This discrepancy has been an ongoing problem which we do not yet fully understand.

5.7. Field reversal and signal processing. – With all the necessary frequencies, intensities and lengths stabilized, one then has to be concerned about reversing the various fields as precisely as possible, without upsetting the servo control systems. The electric-field flip is accomplished by reversing the voltage applied to the electric-field plates. Initially we tried a sinusoidal reversal, but this gave unacceptably large electrical pickup on the detector. We then switched to a square-wave modulation with a few milliseconds of dead time after each reversal before taking data, to allow the transients to die away. The primary problem in obtaining a perfect electric-field reversal is the stray field. There is considerable black magic we have learned for the preparation and handling of the plates which keeps the stray fields to a minimum, typically, a few tens of mV/cm.

For the actual voltage reversal, we have experimented with various solid state and mechanical switches and have obtained the cleanest reversals when we use high-voltage relays. These have the minor annoyance that they are somewhat slow (less than ~ 40 Hz), but the reversal is much more exact than with any solid-state devices we have found. Mercury relays are faster but are limited in the voltage they can handle. To reverse the polarization, we use the same high-voltage relays to flip the voltage applied to the Pockels cell that provides a quarter-wave retardation. In this reversal, the major problem is the birefringence of the Pockels cell which drifts with temperature. However, with careful temperature stabilization this can be reduced to a reasonable level.

The magnetic-field reversal is the easiest; we simply reverse the current flowing through coils using solid-state switches. One of the major concerns when doing any of these reversals is to avoid upsetting any of the servo loop.
This takes considerable care and involves the use of various sample-an
circuits and gates with precise timing to isolate the servos from the trai
We have succeeded in keeping everything stable enough that the nois-
flipping the various fields is as low as when there are no reversals.

The signal processing is the final part of the apparatus, and it is fair-
ple. The current from the photodiode is sent into a very-low-noise curr
voltage converter, as mentioned above. The output voltage of this ampli-
monitored with two different systems. The first is relatively crude (1 : 
10^3), and simply monitors the overall d.c. signal level for normalizatio
second system detects the small changes in the signal due to the PNC n-
tion. First, the signal passes through a low-noise amplifier which subtract
constant voltage so that the d.c. output is close to zero. This near-zero si-
sent into a gated integrator and the signal is integrated during the ti
between each reversal. At the end of each interval the output from the inte-
is digitized and stored in a computer. Each of these numbers is stored v
appropriate label as to the state of E, B, m and polarization. Then the cor-
carries out the next reversal, resets the gated integrator, and the sequ
repeated. Using the offset and gated integrator in this manner, we av
dynamic-range problem encountered in trying to measure a very small m-
tion on top of a large signal.

6. – Systematic errors.

Most of the time spent taking data in these experiments is devoted
study and reduction of potential systematic errors. Our approach to o-
with systematic errors follows the same general analysis used in the ear-
sium and thallium experiments. This procedure starts by considering the
gen general possible case of both d.c. and a.c. electric and magnetic fields
have components in the X, Y and Z directions. Thus we have 12 possi-
components. Next we look at all the combinations of these fields that ca-
rise to a 6S → 7S transition, either electric dipole or magnetic dipole. We
each of these 12 field components to have both flipping and nonflipping (i
forth known as «stray» parts). We then go through the exhaustive list of o-
nations that produce terms that mimic the parity nonconservation by rev-
with all of the possible various reversals. Then we measure the size of the
different field components and, in the process, try to make the stray and
aligned fields as small as possible. We have been able to reduce them to ty-
ly 10^{-4} to 10^{-5} of the main applied field. Using the measured sizes of the c-
ent components, we look at all the vast number of combinations that : 
PNC, and see which ones are significant. The 10^{-4} to 10^{-5} values effect-
mean that any terms involving more than two stray or misaligned compo-
are negligibly small compared to the true (10^{-5} ÷ 10^{-6}) PNC. At the e
Table I. - Potential systematic errors.

<table>
<thead>
<tr>
<th>Systematic contribution (^{(a)})</th>
<th>Range (^{(b)})</th>
<th>Average all data</th>
<th>Daily uncertainty</th>
</tr>
</thead>
<tbody>
<tr>
<td>((\Delta E_y/E)(B_z/B))</td>
<td>(-0.3% \rightarrow +1.1%)</td>
<td>+0.3%</td>
<td>0.4%</td>
</tr>
<tr>
<td>((\Delta E_z/E)(E_y/E))</td>
<td>(-1.3% \rightarrow +0.4%)</td>
<td>-0.1%</td>
<td>0.4%</td>
</tr>
<tr>
<td>((E1M1\xi)(\Delta m = \pm 1))</td>
<td>(-0.8% \rightarrow +4.8%)</td>
<td>+1.7%</td>
<td>0.6% ((\Delta F = -1))</td>
</tr>
<tr>
<td></td>
<td>(-1.1% \rightarrow +6.8%)</td>
<td>+2.4%</td>
<td>0.9% ((\Delta F = +1))</td>
</tr>
<tr>
<td>((\Delta m = 0))</td>
<td>(-0.3% \rightarrow +0.6%)</td>
<td>+0.04%</td>
<td>0.04% ((\Delta F = -1))</td>
</tr>
<tr>
<td></td>
<td>(-1.6% \rightarrow +0.1%)</td>
<td>-0.23%</td>
<td>0.06% ((\Delta F = +1))</td>
</tr>
</tbody>
</table>

\(a\) \(\Delta E_y\) and \(\Delta E_z\) are nonreversing electric-field components, \(B_z\) and \(E_y\) are misaligned magnetic- and electric-field components, and \(\xi\) represents the birefringence of the coating on the output mirror.

\(b\) The range shows largest and smallest daily corrections.

This exercise we found there are three terms that contain two small components, and these are listed in Table I. The first of these terms involves a stray electric field in the Y direction times the (misaligned) magnetic field in the X direction. The second term involves a stray electric field in the Z direction times misaligned component of electric field in the Y direction. And the third term is product of \(E1\) and \(M1\) transition amplitudes times a mirror birefringence factor.

We measure each of the fields and the birefringence involved in these terms while the experiment is running and subtract off their contributions. To do this, we run a set of auxiliary experiments simultaneously, or interleaved with the PNC data acquisition. These auxiliary experiments involve observing the effects on the \(6S \rightarrow 7S\) atomic-transition rate of different hyperfine transition: different laser polarizations and application of additional \(E\) or \(B\) fields. Two points should be emphasized about dealing with systematic errors in this manner. First, it is important to use the atoms themselves so that the same region of space is sampled at nearly the same time as the PNC experiment. Second, these auxiliary experiments must be designed to allow systematic corrections to be measured with an uncertainty that is much less than the statistical uncertainty in the parity nonconservation experiment. It is highly desirable to have a measurement time much shorter than that required to take the parity violation data. If one fails to achieve this, then the uncertainty of an experiment increases because much of the running time is spent in taking data on systematic errors and little on the measurement itself.

In the experiments we have designed, achieving the necessary uncertainties requires a small fraction of the PNC integration time. In Table I, we show the different sizes of the systematic uncertainties for our 1988 experiment, how much they vary from one run to another, and the average correction and uncertainty.
tainty. It can be seen that the typical corrections are a few percent or less, most importantly, the uncertainties in all of these corrections in a given region are less than 1%, and thus much smaller than the statistical uncertainty.

An obvious question is, «Is this analysis foolproof, or did we miss something?». In fact we did miss something; there is no analysis that is at all foolproof. We overlooked a small correction the first time through, although we caught it well before we were ready to publish a result. However, it is essential to discuss the statistical-analysis procedure we used to discover systematic error. This same type of analysis can (and probably should) be used in any precision measurement. It involves using a $\chi^2$ test in a particular track down systematic errors. The Allan variance used to characterize the stability standards is related to this approach.

Our data consist of a large set of numbers, each number corresponding to a current which was integrated for 0.1 s. In the entire data set there are $10^7$ such numbers stored in the computer for analysis. The first step is to find the scatter in the numbers which is due purely to the statistics and has nothing to do with any systematic source. This is accomplished by looking at the fluctuations on the shortest possible time scale where the statistical fluctuations are large. This gives us a standard deviation, $\sigma$, which is most likely purely statistical. In our case, we are doubly sure that it is truly statistical because it corresponds to the shot noise limit for the signal.

Having found $\sigma$, we collect the data into various bin sizes, for example, the first million data points would be one bin, the second million would be two bins, and so on for all the data. This produces 10 bins of data, and we predict how the average values in each bin should distribute based on $\sigma$, the number of points in the bin ($\sigma_{\text{ave}} = \sigma \sqrt{N}$). This hypothetical distribution is compared with the actual distribution. Specifically we find the value of the distribution using $\sigma$ to obtain an uncertainty for the value in each bin. We look up the probability for having that value of $\chi^2$. If the resulting probability is 0.5 or larger, we are confident there was no systematic error that was missed.

Now, by choosing different bin sizes we probe for variations on different time scales. This is quite important because, if the bin size is much larger than the time scale for the variations in some systematic factor, the $\chi^2$ will probably look reasonable. However, when one chooses a bin size that corresponds to the time scale of the variations, suddenly, the probability becomes very low, indicating the presence of some unknown systematic error. The approach does not have to be limited to binning the data by time. It is equally useful to bin it according to any other factor that may lead to some systematic errors. For example, one might also bin the data according to room temperature or to search for temperature-dependent systematic errors.

Of course, this approach only works if the systematic errors vary—il
are always constant you will never see them. However, one can make them visible by changing everything about the experiment that might be important, such as realigning all the optics or replacing critical components. Again, we binned the data corresponding to the different configurations and performed the $\chi^2$ test. This is a remarkably sensitive test for potential systematic errors, and, though it is not generally taught, it is important to keep in mind in any precision experiment. In our case it revealed that we had neglected to consider the $E1M1$ interference correction associated with the off-resonance excitation of the $m$ level. This excitation is forbidden at zero magnetic field but could occur due to the second-order Zeeman effect.

Having carried out all the detailed studies of systematic errors and $\chi^2$ tests we finally achieved the result

$$\frac{\text{Im} \delta_{\text{PNC}}}{} = \begin{cases} -1.639(47)(08) \text{ mV/cm}, & F = 4 \to F' = 3, \\ -1.513(49)(08) \text{ mV/cm}, & F = 3 \to F' = 4, \\ -1.576(34)(08) \text{ mV/cm} & \text{ (average).} \end{cases}$$

The size of the parity-nonconserving mixing is given in terms of the equivalent amount of d.c. electric field that would be necessary to give the same mixing $S$ and $P$ states. As shown, we have measured this mixing for two different hyperfine transitions, the $6S, F = 4 \to 7S, F' = 3$ and the $6S, F = 3 \to 7S, F' = 4$. In both cases, the amount of mixing corresponds to about 1.5 mV/cm. The average of these two is the most important quantity, as we will discuss below. We have measured this to an uncertainty of 2%, which is dominated by the 0.05 mV/cm statistical uncertainty. The systematic uncertainty is about 1/4 of the size of the statistical uncertainty. It should be noted that this systematic uncertainty is different from many systematic uncertainties, in that it is actually the true statistical uncertainty in the evaluation of the systematic correction. Therefore, if the statistical signal-to-noise ratio in the experiment is improved, this uncertainty will be reduced.

In fig. 20, we show a comparison of the different experimental measurements of parity nonconservation in cesium, the most thoroughly measured atom. On top are the two experimental results of the Paris group in '82 and '84 below is our 1985 result, and our 1988 result, with its 2% uncertainty. There is good agreement among all of these numbers. This gives one a certain amount of confidence that no tremendous systematic errors are being overlooked.

In table II, we show a summary of the results from all atomic parity nonconservation experiments. In the first section are the optical-rotation experiments which looked at the 648 nm line of bismuth. These results are somewhat controversial in that the results from the three groups showed substantial discrepancy, as did the theoretical calculations. In retrospect, the former was probably due to systematic errors that were not sufficiently controlled. More recent optical-rotation experiments have shown better consistency, and the uncertainties...
are mostly in the (15 ± 30)% range. The one exception is the recent Oxfo measurement in bismuth which has an uncertainty of only 2%. The Stark-interference experiments are given at the bottom of this table. Most of the cesium measurements we have already mentioned, plus there is thallium result from Berkeley with an uncertainty of 28%.

Before we can consider what these measurements tell us about elem particle physics, we must return to the atomic-structure issue. The that is experimentally measured is the $\delta_{\text{PNC}}$ mixing, which is equal to $Q_5$ the $\gamma_5$ matrix element. This matrix element is found by calculating the

**Table II. – Summary of PNC results.**

<table>
<thead>
<tr>
<th></th>
<th>Experiment</th>
<th>Atomic</th>
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<tbody>
<tr>
<td><strong>Ancient (controversial) history</strong></td>
<td></td>
<td></td>
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<tr>
<td><strong>Modern civilized (?) era:</strong></td>
<td></td>
<td></td>
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<tr>
<td>Pb</td>
<td>(Washington ’83) [16]</td>
<td>±28%</td>
</tr>
<tr>
<td>Bi</td>
<td>(Washington ’81) [18]</td>
<td>±18%</td>
</tr>
<tr>
<td>Tl 1.3 μm</td>
<td>(Oxford ’91) [20]</td>
<td>±2%</td>
</tr>
<tr>
<td></td>
<td>(Oxford ’91) [21]</td>
<td>±15%</td>
</tr>
<tr>
<td><strong>Stark-induced interference:</strong></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Tl</td>
<td>293 nm (Berkeley ’85) [11]</td>
<td>±28%</td>
</tr>
<tr>
<td>Cs</td>
<td>(Paris ’84-’86) [23]</td>
<td>±12%</td>
</tr>
<tr>
<td></td>
<td>(Colorado ’85) [9]</td>
<td>±12%</td>
</tr>
<tr>
<td></td>
<td>(Colorado ’88) [10]</td>
<td>±2%</td>
</tr>
<tr>
<td></td>
<td>all agree with the standard</td>
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structure. Thus, in order to obtain $Q_w$ to 1%, both the experiment and the matrix element calculation must be accurate to better than 1%. As mentioned earlier, the calculation of $\gamma_5$ varies considerably in accuracy from one atom to another. In table II we have given the accuracy quoted for the best calculation for each atom.

Here we will limit our discussion to the cesium atom for which there have been the most abundant and most accurate calculations. Two basic approaches have been employed for these calculations. The first is the semi-empirical method which has been used in Paris, Oxford, Colorado and elsewhere. This approach uses experimental data to determine wave functions which are then used to find the matrix element. This technique is relatively easy. However, it is difficult to make a rigorous evaluation of the accuracy of the calculation, since the relevant experimental data have already been incorporated into the calculation. The estimates for the uncertainty in these calculations are as small as 2%. The second approach is to use ab initio relativistic many-body perturbation theory. The need for accurate cesium PNC calculations has spurred major advances in this field, although the calculations are very long and difficult. The most recent and most accurate results have come from the Novosibirsk group of Flambaum, Sushkov et al., who have achieved a 2% uncertainty, and the Notre Dame group of Blundell, Saperstein and Johnson, who have now reached 1% uncertainty. The advantage of this calculational approach is that there is a fairly clear prescription for evaluating the accuracy of the calculations. The most direct way is to simply use the same calculational technique to determine many properties of the atoms and compare these with experimental data. In this case this means calculating hyperfine splittings, oscillator strengths between man transitions, energy levels and fine-structure splittings for cesium and other alkali atoms. Fortunately, a tremendous amount of experimental data is available for comparison. In all cases, the agreement between the calculations and the experiments has been within 1%. Another technique for estimating the uncertainty in these calculations is to estimate the size of the uncalculated higher-order terms in the perturbation series expansion. This approach also gives an uncertainty of about 1%.

7. - Implications.

In this section, we will consider the implications of the Colorado measurement of PNC in cesium. We will first discuss the significance of the comparison of the two different hyperfine transitions. This difference between the two numbers, $\Delta = 0.126(68)$ mV/cm, is probably not zero. More specifically, this value indicates a 97% probability that $\Delta$ is greater than zero. When we made this measurement, we did not anticipate a nonzero result at this level and therefore, spent a considerable amount of time trying to determine what was
wrong with the data. The result, however, stubbornly persisted. Only lat we discover that an effect of nearly this size had been predicted.

The primary difference between these two transitions is that the n spin is reversed relative to the electron spin. Thus $\Delta$ is a measure of the nu spin-dependent contribution to the PNC signal. Two processes have bee cussed which would cause a nuclear-spin-dependent parity nonconserv. The first is simply the electron-quark portion of the weak neutral current depends on the spin of the quarks. This interaction is characterized by $t_1$ and $C_{2d}$ coefficients. As we mentioned earlier, because of the size of these coefficients and the fact that the effect is proportional to the total nuclear spin not proportional to the number of quarks), this contribution is much sr than the weak-charge contribution. However, it has also been pointed o that there is a substantially larger contribution, called the nuclear anapole moment, which arises from weak interactions within the nucleus. The effect these weak interactions (both charged and neutral) are to mix the parity states of the nucleus, leading to a parity-nonconserving electromagnetic current in the nucleus. This current takes the form of a toroidal helix, and, tfore, has no long-range electric or magnetic fields. Thus it gained the «anapole moment». This phenomenon was first proposed by ZEL'DOVICH in in the general context of parity violation in charged systems[25]. It is not known because people shortly thereafter decided such an effect could new measured. However, because the cesium electrons penetrate the nucleus, spend some time inside the toroidal helix and thereby detect its existence. coupling to the electrons is purely electromagnetic, but, because the under nuclear currents are parity violating, it leads to parity violation in the elec ic transition.

There has been a significant amount of interest in this nuclear anapole ment by the nuclear-physics community and several authors have calculated expected size. The first calculations were by KHRIPLOVICH and FLAMBAUM[33]; their estimates are consistent with our observations. HAXTON et al. [26] also made similar calculations, but have treated the nuclear physics rather ferently. Finally, BOUCHIAT and PIKETTY have done a calculation which is consistent with our result[27]. We have been told by FLAMBAUM that the diences in these calculations are not due to any fundamental difference in the ory, but are a problem of the basic interpretation of nuclear PNC from other experiments. Depending on how one chooses to interpret the other experiment it is possible to obtain very different constants which characterize PNC intations in the nucleus. This emphasizes the need for more accurate data in field. There is hope that these nuclear-anapole-moment measurements can vide these data. The nuclear anapole moment is unique in that it is a PNC tion of the nuclear ground state. The previous measurements on nuc PNC have observed parity mixing of excited, and often rather distorted, clear states where there is considerable uncertainty about the nuclear w
functions. Thus it is clear that future improvements in atomic PNC precision should substantially improve the understanding of nuclear parity nonconservation.

Obviously, the uncertainty due to the nuclear physics is a serious issue in the interpretation of atomic parity nonconservation. If we had measured only a single transition, it would seriously compromise our ability to test the standard model. Fortunately, if we take the average of the measurements on the two hyperfine transitions, as opposed to the difference, the nuclear-spin-dependent part cancels out. In this way, we also cancel out any questions involving the nuclear structure, which is critical in allowing a precision test of the standard model.

From this average and the Notre Dame matrix element calculation, we obtain a weak charge, \( Q_w = 71.0 \pm 2\% \pm 1\% \). If one assumes the standard model is correct, one can then from this extract a value of \( \sin^2 \theta_w \) which is equal to

\[
\sin^2 \theta_w = 0.223 \pm 0.007 \text{ (experimental)} \pm 0.003 \text{ (theoretical)}.
\]

This value of \( \sin^2 \theta_w \) can now be compared with values obtained from other experiments such as the measurement of the \( Z_0 \) mass or the neutrino scattering results, as shown in Table III. In addition to these two measurements, there are many other measurements from high-energy experiments which can be used, but which have lower precision or involve other properties of the \( Z_0 \). We have omitted the latter group because, while these are reputed to be independent measurements, the variations in the values of \( \sin^2 \theta_w \) obtained are much smaller than the quoted uncertainties. This leads to unrealistic \( \chi^2 \) probabilities and suggests that these measurements are not truly independent.

The comparison of the values of \( \sin^2 \theta_w \) provides a precise test of the standard model. It is worth noting that in other tests of the standard model, particularly those involving the comparison of neutrino scattering and \( Z_0 \) data, the uncertainty in the mass of the top quark introduces an uncertainty of 0.003 in the relative values of \( \sin^2 \theta_w \)[2]. However, the comparison of the atomic and the \( Z_0 \) mass values is unique in that dependence on the top quark mass is essentially identical in the two cases. Because the atom is sensitive to a different set of electron-quark couplings and a different energy scale, this comparison, however, is very sensitive to new physics which is not contained in the standard

<table>
<thead>
<tr>
<th>Table III. - ( \sin^2 \theta_w ) values.</th>
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<tr>
<td>( \delta_{\text{PNC}} ) (experiment) + ( \langle \gamma_5 \rangle ) (theory)</td>
</tr>
<tr>
<td>( \Rightarrow Q_w = 71.0 \pm 2% \pm 1% )</td>
</tr>
<tr>
<td>( \Rightarrow \sin^2 \theta_w = 0.223 \pm 0.007 \pm 0.003 )</td>
</tr>
<tr>
<td>( W ) mass ( \Rightarrow 0.2320 \pm 0.0007 )</td>
</tr>
<tr>
<td>neutrino scattering ( \Rightarrow 0.233 \pm 0.003 \pm 0.005 )</td>
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</table>
model. Proposed examples of such new physics include technicolor an bosons, which occur in many models.

Figure 21 shows the values of $C_{1d}$ vs. $C_{1u}$ as determined by a mode
dent analysis of the experiments, along with the standard model. The hatched area is the constraint from the SLAC deep inelastic scattering experiments. In contrast, the constraint set by our cesi
ment and the Notre Dame theory is the narrow solid line which is n thogonal and, in particular, constrains the value of $C_{1d}$ much more se
crossed line shows the values allowed by the $SU_2 \times U_1$ Weinberg
/Glashow theory. The point on that line is determined by the value c sin$^2 \theta_w$. This figure shows quite clearly that any new physics that woul
change in the value of $C_{1d}$, but not affect the other coupling constan

![Diagram](image-url)

**Fig. 21.** - Constraints on the $C_{1d}$ and $C_{1u}$ coupling constants by experimental ments. The hatched region is from SLAC deep-inelastic-scattering data, while line is from atomic PNC. The $SU_2 \times U_1$ line is the standard model value as a fu sin$^2 \theta_w$. 
only be revealed by the atomic-parity-nonconservation measurements. The arrows on this figure show how a few popular proposed models would shift the values of these two coupling constants to a different place in the plane. Because the current atomic-physics line passes through the standard model point, there is no indication of the existence of new physics. However, this does put constraints, in some cases quite severe, on the parameters of models that propose such new physics.

One example which has drawn considerable attention in the last few years is how atomic PNC results constrain the proposed mechanism known as technicolor, or, more generally, dynamical symmetry breaking involving heavy particles. This type of new physics has been characterized in terms of $S$ and $T$ parameters which enter directly into $Q_w$ [28]. Generic technicolor models predict the $S$ parameters should be around $\pm 2$ or somewhat larger [28]. From the comparison of results just mentioned, one finds that the atomic PNC yields a value of $S$ which is $-2.7 \pm 2 \pm 1.1$ as given in ref. [29]. Thus one finds that technicolor is on somewhat shaky ground, although the atomic PNC experimental results are not good enough to completely rule it out.

Finally, atomic PNC provides the best constraints on many models that involve additional neutral $Z$ bosons. While there are many papers on this subject (see references in ref. [2]), we note particularly the paper by Mahanthappa and Mohapatra [29]. They consider 11 different models with additional $Z$'s which have been proposed, and they find that, in 8 of these 11 cases, atomic PNC provides the most severe constraints. Thus it is clear that the cesium PNC results are providing information on elementary-particle interactions that is not available from any other source at the present time.

8. – Future improvements.

8'1. Near term. – While atomic PNC experiments are providing useful information, it is clear that more precise results would be desirable and useful. The mass of the $Z_0$ is now known to around 1 part in $10^8$. If atomic PNC results could be improved to that level, we would have a 10-fold improvement in the test of the standard model and correspondingly improved sensitivity to possible new physics. With this in mind, we would like to discuss our efforts to improve the cesium PNC results. Work is also under way to improve atomic-parity-nonconservation measurements by several other groups: In Paris, the Bouchiat group is building a new experiment that involves stimulated-emission probing of the excited state in cesium. At Oxford and Washington, experiments are under way to obtain more precise optical-rotation measurements in thallium. At Berkeley efforts continue to obtain a more precise Stark interference measurement in thallium. All of these experiments have been under development for a number of years, and we hope to have results in the not too distant future.
In our efforts to improve the Colorado 1988 experiment, our primary focus has been on improving the signal-to-noise ratio. This is clearly the major concern of our experiment since the statistical uncertainty was much larger than the systematic uncertainties. We also took into account the fact that the scattered laser light was a major nuisance requiring frequent readjustment of the laser optics, and the transparent conducting coatings used to protect the cesium vapor from exposure to the cesium vapor, and lead to interruptions in the experiment while they were replaced. With these in mind we built a new apparatus that uses an optically pumped atom which, in principle, should provide 16 times more atoms since there are 16 m levels of the 6S state. We now have better mirrors for our pow-up cavity; these increase the build-up by about a factor of 10, resulting in a build-up of 15000. A third improvement is using downstream detectors to detect 6S → 7S excitation. This concept is illustrated in fig. 22, which shows a schematic of the new apparatus.

After leaving the oven, the cesium atomic beam is optically pumped into the single F and m_F level by light from two diode lasers which drive two hyperfine transitions of the 6S → 6P_{3/2} transition. The atoms in the single m level propagate down the atomic beam and intersect the power build-up cavity where they are excited to the 7S state. They then have a 70% probability of caying back down into the 6S hyperfine level which was previously depleted. The atoms continue down the optical beam in this state until they reach the probe region. In this region, light from another diode laser again excites them to the 7S state.

---

Fig. 22. — Schematic of the new Boulder cesium PNC apparatus.
6S → 6P_{3/2} transition. However, here we excite a cycling transition (F = 4 → F' = 5 or F = 3 → F' = 2). On a cycling transition the atom returns only to same initial state and hence can be excited many times. Typically 1000 i.r. photons are scattered for each 6S → 7S excitation. We detect this fluorescence determine the 6S → 7S excitation rate. This detection scheme provides a substantial amount of amplification, yielding a detection of about 200 photons a 6S → 7S transition, instead of the 0.3 detected in the previous apparatus.

In addition, since the detection takes place at a different region from the citation region, we can now construct our electric-field plates out of any material. This greatly simplifies their construction and increases their longevity. Also, scattered light from the green laser light is now negligible, as is detection noise, because the signal size is much larger. All of these improvements would suggest that the experiment should be much easier. In fact, there have been major headaches and delays with this approach, and it is educational to consider what has gone wrong and what lessons can be learned about doing experiments at the frontier of laser spectroscopy.

We will now discuss the unexpected problems we encountered in making this «improved» experiment work. The first problem was noise in the optical pumping and resonance fluorescence detection regions due to the fact that we were using diode lasers. Diode lasers have very rapid (ns) fluctuations in the optical phase. Through a somewhat obscure process, this leads to very-low-frequency fluctuations in the atomic-transition rate. This was quite puzzling when we first observed it, and has now been explained in a series of papers by Zoller and collaborators[30]. While this has become interesting atomic/optical physics to a number of people, to us it is a major experimental problem. To avoid this problem one must have a feedback system capable of providing gigahertz bandwidth correction signals in order to eliminate the noise at (10 ∆ 20) Hz, when we detect our signals. In the first attempt we used optical feedback from narrow-band resonators, as demonstrated by Hollberg and co-workers[31]. This approach gave low-noise signals, but the locking was not reliable enough to allow three diode lasers to operate for reasonable periods of time. After considerable additional work, we settled on using optical feedback from diffraction gratings[32]. These gave much more reliable performance, but the noise levels were still unacceptably high. We solved this problem by finding a laser manufacturer whose instruments gave superior performance when operated with grating feedback. Thus we have finally succeeded in producing a very reliable source of diode laser light which provides very good signal-to-noise ratio in excitation of narrow-band atomic transitions. Specifically, we can now achieve a noise-to-signal ratio of 3 · 10^{-6}/Hz^{1/2} when exciting a 10 MHz wide atomic transition. The necessary laser has a combination of optical, mechanical and current feedback, as shown in fig. 23. The grating which provides the optical feedback is mounted on a piezoelectric transducer which allows mechanical adjustment of the grating position. This holds constant the length, and hence the frequency
of the optical cavity. For faster corrections to the cavity resonant frequency, the laser current to achieve the highest level of stabilization.

The second major problem in the new experiment was background in the supposedly empty $F$ state. It is relatively easy to deplete one $F$ level by low-intensity atomic beam very well ($< 10^{-4}$). However, with a more intense beam there are a number of mechanisms that can repopulate the empty $F$ state. For example, collisions between atoms in the beam and surfaces or other atoms (particularly oxygen) is the first mechanism. We have eliminated this source by improving the vacuum and carefully positioning the collimating surfaces.

Another contribution to the background, which appears to be from atoms in the wrong $F$ state, actually comes from the excitation of the other hyperfine levels. The tail of the spectral distribution of light in the probe beam. We have eliminated this source by sending the probe light through an interferometric cavity which blocks out the tails of the spectral distribution. The third and most serious source of atoms in the wrong $F$ state has been the multiple scattering of the optical-pumping light. The optical-pumping process produces fluorescence, which can travel down the atomic beam and re-excite the atoms out of the $F$ state. The number of atoms pumped back into the empty state scales with the square of the atomic-beam intensity. We have found several ways to reduce the background: multiple pumping beams (the «clean-up» beam in fig. 22), the optical-pumping transition which minimizes the scattered fluorescence.
using the photon blocking collimator. This is a collimator with very thin blades which allows only the highly collimated photons to pass through. Finally, even with all these steps, the background was still too large, and we ultimately had to reduce our atomic-beam intensity. In spite of all these setbacks and delays, this new, improved experiment is now operational and we are taking data with a signal-to-noise ratio several times better than that of the 19 experiment.

A very painful lesson has been brought home to us in carrying out this improved experiment. When one is probing a region of technology and physics which is unexplored, it is important to step warily, and to keep all your options open. In terms of an experiment, this means you should keep the apparatus flexible and be ready to adapt, as we mentioned earlier. In this experiment we were somewhat seduced by the fact that this approach seemed to solve all old problems, and we committed ourselves to a design that turned out to be filled with major unexpected difficulties.

8.2. Long term. – As the experimental accuracy improves beyond 1%, the principal limitation on the usefulness of atomic PNC will become the atomic theory. There have been credible speculations that it will be possible to calculate the theory in cesium to a part in \(10^3\). However, it is not clear when the calculations will be completed, and the question of how to check their accuracy becomes a major issue.

We have begun a longer-term experimental project that avoids the atomic theory question. The basic idea is to compare precise measurements of atomic PNC for different isotopes of cesium. The weak charge is sensitive to the number of neutrons, and hence will change for different isotopes. The atomic matrix element, however, depends on the electronic structure and is almost independent of the number of neutrons. If one then looks at appropriate combinations and ratios of experimental results, for example

\[
\frac{\delta C_{\text{PNC}}^{130} - \delta C_{\text{PNC}}^{150}}{\delta C_{\text{PNC}}^{130} + \delta C_{\text{PNC}}^{150}} = \frac{Q_{\text{w}}^{130} - Q_{\text{w}}^{150}}{Q_{\text{w}}^{130} + Q_{\text{w}}^{150}},
\]

the atomic matrix element will drop out, leaving a ratio of weak charge which can be directly compared with standard model predictions. In this manner we hope to achieve measurements that can be compared with the standard model predictions at the part in \(10^3\) level. There are two major obstacles to carrying out these experiments. First is the need for even better signal-to-noise ratios. Second, and most critical, is the need to carry out PNC measurements with small atomic samples, rather than the many grams used in the atomic-beam measurements. This requirement is necessary because all the other isotopes of cesium are radioactive and can only be obtained and used in small quantities. We propose to overcome both of these obstacles by using the new technology.
laser trapping. This will improve the signal-to-noise ratio because it is even easy, to obtain optical thicknesses in trapped-atom samples 10 or 10 larger than can be achieved in our atomic beam.

It is more difficult to show that optical trapping will allow the expe to be done with very small atomic samples ($10^{10}$ atoms). We are currentl ing on this problem. The approach we are using starts with a very small of a given isotope (short-lived isotopes will be produced at an accelerato: longer-lived isotopes can be brought to our laboratory), which is injected special cell where the atoms will be efficiently captured by a laser trap ()
We have carried out detailed studies on capturing atoms from a vapor: are currently developing wall coatings which will allow the cesium at bounce around inside the cell without sticking until they are captured. Finary work with silane coatings has been quite encouraging.

Once the atoms are captured, PNC measurements can then be carried the cold dense samples. If all goes according to plan, the next decade v high-precision measurements of PNC in a number of cesium isotopes. Ti provide detailed information on the nuclear anapole moment and a very test of the standard model.

***

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REFERENCES


