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**NEUTRAL CURRENTS AND PARITY BREAKDOWN
IN ATOMIC TRANSITIONS:
THREE PROPOSED EXPERIMENTS**

S. D. Bloom

February 8, 1975

MASTER

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Abstract

This paper describes three proposed experiments for observing the breakdown of parity in atomic transitions due to the exchange of neutral, parity-violating currents arising from some of the new gauge models (e.g., the Weinberg model) for the weak interaction. The experiments are based on exploiting a suggestion, by Bouchiat and Bouchiat, that modern laser technology be utilized to produce intense, monochromatic, and polarized photon beams with which to excite forbidden atomic transitions of the basic form $\ln s^{1/2} \rightarrow \ln s^{1/2}$. The asymmetries (of the order of 10^{-6}) in the de-excitation processes then signal the presence of the parity-violating components due to

the neutral currents. In all three experiments suggested here, the use of multiple (uncollimated) atomic beams as targets forms a basic part, and their advantages over a temperature-equilibrium vapor are described. The first experiment uses ^{55}Cs atomic beams as a target; the second uses ^{85}Rb in conjunction with a super-strong magnetic field (~ 80 kG); the third uses ^{81}Tl and requires frequency doubling of the exciting laser beam. All three experiments appear to be quite feasible, and, given the requisite equipment (much of which is or soon will be commercially available), they could yield definitive results in a period of a few months.

I. Introduction

In this section, we present an elementary review of the present picture of the structure of the electromagnetic and weak (or β) interactions.

This picture (or, more properly, theory) is undergoing a drastic revision right now. In some respects, in fact, it is changing almost on a week-to-week basis. An excellent layman's review of the subject has been given by Weinberg in a recent issue of *Scientific American*,¹ in which he emphasizes this rapidly changing aspect of the subject, among others.

Some features of Weinberg's approach will be retained here, but our main object is to present an understandable basis for the use of atomic transitions to study the weak interactions as they affect the polarization and parity properties of certain atomic transitions. This approach to the study of the weak interactions was not among those considered in Weinberg's papers¹⁻⁴ (see especially Ref. 4), but it was explicitly studied by Bouchiat and Bouchiat,^{5,6} who found, rather surprisingly, that combining modern laser technology with circular polarization could in certain cases provide a very sensitive test for the presence of the so-called weak-interaction neutral currents. These neutral currents do not conserve parity and are an important feature of the new unified theory of weak and electromagnetic interactions,⁷⁻¹⁰ at least in the simplest or so-called Weinberg version.⁴

Of course, the exchange of either charged or neutral currents among the "strongly" interacting particles

(for example, the exchange of π^+ or π^0 mesons between two nucleons) has been familiar for decades. However, the "strong" or nuclear interaction is strictly parity-conserving, as far as is known. In fact, the simplest version of the nucleon-pion interaction, the so-called strong pseudoscalar-pseudoscalar interaction (we call this H_s) is worth introducing at this point as an illustration of a very simple fermion-boson coupling:

$$H_s \approx i(4\pi)^{1/2} \bar{\psi}_N \gamma_5 \vec{\phi}, \quad (1a)$$

where

$$\vec{T}_N \equiv \text{nucleon isospin operator } (t_1, t_2, t_3), \quad (1b)$$

$$\gamma_5 \equiv \text{nucleon Dirac "pseudoscalar" operator}, \quad (1c)$$

$$\vec{\phi} \equiv \text{the isovector pion wave function}, \quad (1d)$$

$$\vec{\phi} \equiv \phi_+ \vec{T}_+ + \phi_0 \vec{T}_0 + \phi_- \vec{T}_-, \quad (1e)$$

where ϕ_+ and ϕ_- are the wave functions for charged and neutral pions, and the $\vec{\phi}$ indicates a vector in isospin space, as indicated by Eq. (1e). The coupling strength is represented by f , a positive-definite number. Note that $\vec{\phi} = \vec{\phi}(\vec{r}, t)$; i.e., it is a function of position

and time as well as being an isovector. The interaction H_S leads, for example, to the following reaction:

$$p \rightarrow \pi^+ + n \equiv \langle \pi^+ n | H_S | p \rangle. \quad (2a)$$

$$\langle \pi^+ n | H_S | p \rangle \equiv \text{---} p \text{---} \begin{array}{c} \nearrow n \\ \searrow \pi^+ \end{array} \quad (2b)$$

where we have taken some pains to define the basic equivalence of the reaction leading from p to $(\pi^+ n)$, the matrix element for it, and its Feynman diagram. Since there is only one vertex, this is an example of a first-order or "contact" process. It will serve as the prototype for much of what now follows and is particularly interesting in that it really can be observed.

Of course, the nuclear interaction* is only one of the four known ways that particles can interact. Of the other three, two are characterized by the exchange of only neutral bosons or "currents" (electromagnetism and gravity), and, up until now, one has been characterized by the effective exchange of only charged currents (the weak or β interaction). This curious asymmetry and others (like the unique breakdown of parity in the weak interaction) were attacked in part by Weinberg in 1967.² The first two sentences from that paper are of interest here: "Leptons interact only with photons, and with intermediate bosons that presumably mediate weak interactions. What could be more natural than to unite these spin-one bosons into a multiplet of gauge fields?" In somewhat more mundane terms, what Weinberg is suggesting is that the time to unify the weak and electromagnetic theories may now have arrived. The suggestion is not altogether new, and Weinberg makes reference to some earlier efforts.^{2,4} However the "gauge" approach to the problem, plus certain other requirements imposed by Weinberg (and also by Salam⁵), is new and does lead to certain conclusions and concomitant effects having to do in particular with the exchange of neutral but parity-nonconserving spin-one bosons. Some experimental tests have been made, and more are in progress.

*Using the word "interaction" might lead to some confusion, since this often is interpreted as meaning the interaction between two nucleons. This would require the addition of another nucleon line to the right-hand side of Eq. (2b), which would then lead to another vertex or matrix element. This second-order process is what is meant by a charge-exchange reaction. A similar diagram is given in Fig. 1b.

The initial results are very encouraging, and some account of the status to date can be found in several references quoted by Weinberg.⁶

We will not pretend to be up-to-date on this point. In any case, the amount of published material on the subject is already large and is increasing rapidly, including many review articles.⁶ However, we might add that Weinberg's review at the recent Coral Gables Conference¹⁰ is an excellent (though far more technical) second step after Ref. 1.

To put the simplest elements of the results of the "simple" Weinberg theory in focus, the most effective way by far is to use the Feynman diagrams for the three basic scattering reactions issuing from the unified approach. This is exactly what Weinberg himself does in Ref. 1; however, we will generalize his diagrams somewhat and add the matrix elements corresponding to each. We will limit ourselves only to such diagrams and interactions as pertain to atomic structure.

Figure 1 is thus the starting point of our discussion. Besides the usual "bra-ket" notation, we note the following definitions:

$$p, n \equiv \text{protons, neutrons; } N \equiv \text{nucleon.} \quad (3a)$$

$$e, \nu \equiv \text{electrons, neutrinos.} \quad (3b)$$

$$A_\mu(\vec{r}) \equiv \text{the ordinary relativistic vector-potential.} \quad (3c)$$

$$\mu \equiv 1, 2, 3, 4. \quad (3d)$$

$$j_\mu(\vec{r}) \equiv \text{the neutron four-vector current operator.} \quad (3e)$$

$$W_\mu^\pm(\vec{r}) \equiv \text{the relativistic potential for the doubly charged } (\pm) \text{ spin-one meson. It mediates ordinary } \beta\text{-decay, and its mass is (currently) } \geq 20 \text{ GeV. Its gyromagnetic ratio is 2, like the electron's.} \quad (3f)$$

$$Z_\mu(\vec{r}) \equiv \text{the relativistic vector potential for the neutral spin-one meson. Its mass is } \geq 24 \text{ GeV.} \quad (3g)$$

$$\gamma_5 \equiv \text{the standard pseudoscalar Dirac operator} \\ \equiv \gamma_1 \gamma_2 \gamma_3 \gamma_4. \quad (3h)$$

Fig. 1. The three fundamental exchange particles involved in electromagnetic (γ) and weak ($W^\pm$, Z) interactions. The γ is neutral, the $W^\pm$ carries $\pm$ charges, and the Z is neutral. These last three particles are predicted by the Weinberg "simple-version" unified theory. The mass of the $W^\pm$ is ≥ 24 GeV and of the Z ≥ 20 GeV.

Note that all mesons have spin one (just like the photon) and all interactions in Fig. 1 are contact interactions; i.e., everything for each half of the three diagrams occurs at one given point ($\vec{r}, t$). For example, in Eq. (2b), which is the prototype of each half of the three diagrams, a matrix element would have the following (time-dependent) form, in more common terminology:

$$\langle \alpha' n | H_{\vec{r}}(p) \equiv \text{const.} \int \psi_n^\dagger(\phi^*)^* \gamma_5 \psi_p d^3 \vec{r}. \quad (4)$$

where the γ_5 is a pseudoscalar operator which effectively changes sign when $\vec{r} \rightarrow -\vec{r}$, the parity operation. (Remember, pions are pseudoscalar particles, and p and n are identical particles or "states," except for their isospin projection t_3 .)

One important difference between Fig. 1a (the electromagnetic interaction) and Figs. 1b and 1c (the unified weak interaction) must be heavily emphasized at the outset, and that is that the factor $(1 + \gamma_5)$ appears in the latter but does not appear in the former. This means that parity breaks down completely for the weak interaction but is conserved completely for the electromagnetic interaction. In particle physics,

the weak interaction is called a "chiral" theory (see below).

How does all this (and much more - see Ref. 8) come about? We will avoid all details here; however, we will note the following three characteristics of the unified theory, with the prefatory remark that these characteristics were designed into the theory from the start:

- (1) The Lagrangian or the Hamiltonian, including (almost) all interactions, is built to be gauge invariant as well as Lorentz invariant (covariant).
- (2) It is also built to be "manifestly" charge and mass renormalizable.
- (3) Boson masses, isospin breakdown for the relevant particles, etc., are created through the judicious (although "spontaneous") introduction of a nonzero vacuum expectation value (VEV) for one component of a new scalar isodoublet boson, called ϕ by Weinberg.^{2,6} (We will have more to say about this extra and very elusive particle later; notice that it does not show in Fig. 1.)

It is because of point (1) that this entire approach is generally called a "gauge" formulation.

These three points are not simple, nor are they unconnected. Furthermore, all three taken together do not give rise to a unique (though unified) theory of weak and electromagnetic interactions. It is clear that other restrictions and maybe other quantum numbers are necessary, and Weinberg discusses these points in fair detail.⁶ Nonetheless, if we accept (at least for the moment) the simplest version of the Weinberg model,⁴ the number of unknown parameters is reduced to one: the Weinberg angle, θ_w (see below). However, the Weinberg model has several problems yet unsolved, not the least of which is its failure (for the moment) to deal with the strange particles ($S \neq 0$). Protons, neutrons, and pions have $S = 0$; i.e., they are "ordinary" or at least nonstrange. Herein we will ignore the coupling to all strange particles.

The spontaneous breakdown of the VEV for the ϕ meson is a manifestation of the so-called Goldstone bosons. These are usually massless and were predicted long ago to be a problem in any gauge-invariant theory with nonzero VEV's for total energies^{7,8,11} for such particles. The question is, do they really exist? We will not go into this very complex question but note that Ref. 7 discusses the so-called "Higgs loophole," which evades the Goldstone-boson problem (almost) completely. The only vestige of it left in the simple Weinberg theory is the isoscalar-scalar meson ϕ_1 , which itself has (for the moment) a practically indeterminable mass (called M_1). It is the nonvanishing value for the VEV of the doublet ϕ which leads to the breakdown of the isosymmetry between electron and neutrino, so that the electron acquires a mass and an electromagnetic charge (i.e., coupling), the neutrino is decoupled from the electromagnetic field, and the isospin breaks down rather completely for the electron-neutrino system. The ϕ_1 boson not only may have an extremely high mass (> 1000 GeV?) but also has a coupling to electrons about 5000 times weaker than the now (more-or-less) universal coupling strength $(e^2/\hbar c)^{1/2} = 0.085$. Thus we will ignore it (as Weinberg does⁴) in all that follows.

We are now ready to write down the results of the unified theory (for atomic purposes) in terms of two interaction Hamiltonians: 1) a Hamiltonian connecting the ordinary photons and the new bosons ($W^\pm, Z_0$) with electrons, neutrinos, and nucleons,* which we will call h_b , and 2) a Hamiltonian of the usual low-energy current-current kind connecting electrons, neutrons, and nucleons to each other, which we will

call h_f . As noted, we will ignore strange particles ($S > 0$); the hope is that they will not affect our results significantly. Pions we cannot ignore, however; but we will return to this point later. Also, we should note that we do not know if the neutral vector currents are conserved, as the charged vector currents are. This question is dealt with in detail by Pais and Treiman¹²

With these reservations, we now give the two Hamiltonians h_b and h_f [see Eq. (3)]:

$$\begin{aligned} h_b = & \frac{ig}{2\sqrt{2}} \left\{ \sigma_{ce}^\lambda + \sigma_{cN}^\lambda \right\} W_\lambda^\dagger(\vec{r}) + \text{h.c.} \\ & + \frac{ig}{2 \cos \theta_w} \left\{ \sigma_{nN}^\lambda - 2 \sin^2 \theta_w J_N^\lambda + \sin^2 \theta_w \bar{J}_N^\lambda \right. \\ & + \cos^2 \theta_w \frac{(\tan^2 \theta_w - 1)}{2} \bar{J}_{nc}^\lambda \left. \right\} Z_\lambda(\vec{r}) \\ & + ic \left\{ J_e^\lambda + J_N^\lambda \right\} A_\lambda(\vec{r}), \end{aligned} \quad (5a)$$

where

$$g'/g \equiv \tan \theta_w, \quad (5b)$$

$$\begin{aligned} c &= \frac{g g'}{[g^2 + (g')^2]^{1/2}} \approx g \cos \theta_w \equiv \left(\frac{1}{137} \right)^{1/2} \\ &= 0.085, \end{aligned} \quad (5c)$$

$$\sigma_{ce}^\lambda = \bar{\psi}_e \gamma^\lambda (1 + \gamma_5) \psi_e, \quad (5d)$$

$$\sigma_{ne}^\lambda = \bar{\psi}_e \gamma^\lambda (1 + \gamma_5) \psi_e; \bar{\sigma}_{ne}^\lambda = \bar{\psi}_e \gamma^\lambda (1 - \gamma_5) \psi_e, \quad (5e)$$

$$J_e^\lambda = \bar{\psi}_e \gamma^\lambda \psi_e; \bar{J}_e^\lambda = \bar{\psi}_e \gamma^\lambda \gamma_5 \psi_e, \quad (5f)$$

$$\sigma_{cN}^\lambda = \bar{N} \gamma^\lambda (1 + \gamma_5) N, \quad (5g)$$

$$\sigma_{nN}^\lambda = \bar{N} \gamma^\lambda (1 + \gamma_5) \gamma_5 N, \quad (5h)$$

$$J_N^\lambda = \bar{N} \gamma^\lambda \left(\frac{1}{2} + \gamma_5 \right) N. \quad (5i)$$

*The exclusion of muons is for simplicity only. To include muons, one simply adds a muon term to each electron term. This remark applies to the two kinds of neutrinos as well. The simple theory (at this stage) is completely symmetrical in the exchange of the two kinds of leptons.

Normalized units ($\hbar = c = 1$) have been used throughout Eq. (5). Notice that we have three types of currents (basically) for the electrons and the nucleons: two "chiral" currents of the type $\mathcal{J}_{\text{ex}}^\lambda$ and $\mathcal{J}_{\text{nx}}^\lambda$, where e stands for charge exchange and n stands for "neutral" or no-charge exchange, and the usual electromagnetic current J_X^λ . The symbol x stands for either e (electron) or N (nucleon). The "antichiral" current $\mathcal{J}_{\text{ne}}^\lambda$ is a unique feature of the neutral electron current.

$$h_1 = -\frac{G_\beta}{\sqrt{2}} |\mathcal{J}_{\text{ce}}^\lambda + \mathcal{J}_{\text{cn}}^\lambda|^2 - G_\beta \sin^2 \theta_w \mathcal{J}_{\text{ne}}^\lambda + \frac{1}{2} (\sin^2 \theta_w - \cos^2 \theta_w) \mathcal{J}_{\text{ne}}^\lambda + \mathcal{J}_{\text{nn}}^\lambda - 2 \sin^2 \theta_w J_N^\lambda, \quad (6a)$$

where

$$\frac{G_\beta}{\sqrt{2}} \equiv \frac{g^2}{8m_W^2} = 8.8 \times 10^{-44} \text{ MeV-cm}^3, \quad (6b)$$

$$m_W \equiv \text{mass of the } W^\pm \text{ boson } (\gtrsim 20 \text{ GeV}). \quad (6c)$$

Summation over the indices $\lambda = 1-4$ is implied in both Eqs. (5) and (6), and pions have been ignored. Equation (6a) is the low-energy limit, giving rise to the customary current-current version of the β -interaction. For completeness we note that this requires only that the electron and nucleon energies be small compared to m_W ($> 20 \text{ GeV}$) and m_Z ($> 24 \text{ GeV}$). We have left out the electromagnetic interaction, which has the usual low-energy limit of $e^2/(r_{12})$. In fact, our concern will only be with the second and third lines of Eq. (6a), since these are the only ones which could lead to observable parity-violating effects in atomic transitions, as we shall show. The first line of Eq. (6a) leads to the usual type of β -decay, and it can (in second-order) lead to parity-violating effects, but too small to be observable, as we shall also show.

We have said nothing about the actual magnitudes of g and g' , the "new" coupling constants, and, since our main concern is really with Eq. (6a), we really need not. This is because G_β , the β -coupling constant, is well known from many experiments.¹³ However,

we might note that best values for the present⁸ are, in normalized units ($\hbar = c = 1$),

$$\begin{aligned} e &= 0.085, \\ g &= 0.147, \\ g' &= 0.104, \\ \sin^2 \theta_w &= 1/3, \\ g'/g &\equiv \tan \theta_w. \end{aligned} \quad (7)$$

Note how close e , g , and g' are, which lends a good deal of plausibility to the unified character of at least this form of the unified approach.

Both Eqs. (5a) and (6a) are really simpler than they may appear, and Eqs. (1), (2), and (7) provide sufficient guidance to understand their import rather completely. The interaction between the bosons and fermions (h_b) has the form of a total current (although we have not included all the possible terms) coupled to a vector potential, which last has the same form whether it is an electromagnetic boson $\{\vec{A}_\mu(\vec{r})\}$, $\mu = 1-4$, which is massless, or the charged "weak" interaction boson $\{W_\mu^\pm(\vec{r})\}$, which has a relatively gigantic mass ($> 20 \text{ GeV}$). The chiral ($\mathcal{J}$) or ordinary (J) currents are indeed currents. For example, $J_\mu^a(\vec{r})$ is the μ th component (for convenience we suppress the dependence on t) of $\vec{J}_a(\vec{r})$:

$$\vec{J}_a(\vec{r}) \equiv [J_{ex}(\vec{r}), J_{ey}(\vec{r}), J_{ez}(\vec{r}), i\epsilon_{abc}(\vec{r})], \quad (8)$$

where the slash notation $\vec{J}$ means a four-vector. The fourth component is the familiar charge-density $\rho_a(\vec{r})$, except for the factor ic . In the Dirac formalism this is simply γ_4 , in operator form. A more formal definition in the Dirac notation has already been given in Eq. (5f), but Eq. (8) emphasizes the dependence on $\vec{r}$ (and t , if necessary), that is to say, the "ordinariness" of what is meant by a current. In fact $\square \cdot \vec{J}_a = 0$, which is to say that the four-dimension divergence of the electromagnetic current vanishes, as is well known. The chiral currents $\mathcal{J}$, on the other hand, which are distinguished by the (positive) chiral operator $(1 + \gamma_5)$, are not divergenceless because of the presence of the γ_5 or pseudoscalar term. It is the $(1 + \gamma_5)$ operator which corresponds to the V-A or complete-parity-breakdown interaction form used 16 years ago by Feynman and Gell-Mann.¹⁴ It should be no surprise to find it in its proper place in the unified theory, only now we have the neutral currents coupled to the

neutral spin-one bosons $[Z_\mu(\vec{r})]$ as well as the long-ago-surmised charge-exchange currents coupled to the charged spin-one bosons $[W_\pm(\vec{r})]$. However, we must not pass without noting that in fact Feynman and Gell-Mann did speculate on the possibility of neutral currents in their 1958 paper.¹⁴

To put Eqs. (5a) and (6a) in more accessible and hopefully more comprehensible form, we will immediately specialize to the atomic case, which is the principal concern here. In an atom we have only to deal with nucleons and electrons; all other currents can be ignored.* Equations (5a) and (6a) reduce very simply to [see Eqs. (3b-i) for definitions]:

$$h_b(\text{atomic}) = \frac{ig}{2 \cos \theta_w} \left\{ g_{nN}^\lambda - 2 \sin^2 \theta_w g_{wN}^\lambda + \sin^2 \theta_w g_{ne}^\lambda + \cos^2 \theta_w \frac{(\tan^2 \theta_w - 1)}{2} g_{ne}^\lambda \right\} \times Z_\lambda(\vec{r}). \quad (9a)$$

and, dropping the subscripts n ,

$$h_i(\text{atomic}) = -\frac{G_F}{\sqrt{2}} \left\{ \frac{1}{2} (1 - 4 \sin^2 \theta_w) \bar{J}_e^\lambda J_p^\lambda + \frac{1}{2} \bar{J}_e^\lambda J_n^\lambda + \frac{1}{2} (1 - 4 \sin^2 \theta_w) \left[\bar{J}_e^\lambda J_e^\lambda + \left(\sum_{k=1}^3 (\sigma_P^k - \sigma_n^k) J_e^k \right) \right] \right\}. \quad (9b)$$

where [in Eq. (9b)] (σ_P^k, σ_n^k) are ordinary three-dimensional nucleon-spin operators. Also, we have left out all but the parity-violating terms coupling electrons to nucleons and electrons to electrons. Unfortunately, we have had to add the antichiral current $\bar{J}_{ne}^\lambda$ and the pseudo or axial vector current $\bar{J}_e^\lambda$. But they are simple changes from their counterparts, as

*This statement need not be true, strictly speaking. The nucleus spends a significant part of its time as a pion-nucleon mix. However, if the neutral vector current is conserved, then, just as in electromagnetism and the weak (vector) charge-exchange interaction, the virtual pion states make no difference (at the low energies characteristic of atoms). We would get the same answer just by pretending we have nothing but neutrons and protons to deal with. Unfortunately we do not know if the neutral vector currents are conserved or not (see Ref. 13).

noted in Eqs. (5e,f). Needless to say, it is exactly the pseudoscalar-scalar products like $J_e^\lambda J_e^\lambda$ which are parity-violating. Also, we note that we have purposely omitted the parity-violating nucleon-nucleon term; because it cannot affect atomic structure. It does, however, introduce parity-violating effects in nuclear structure, which could change the picture for the expected parity violation in nuclear states. In all considerations so far, only the charge-exchange currents have been considered for this case.¹⁵

It is important to recall that the "ordinary" or charge-exchange β interaction would introduce parity-violating components into atomic structure even in the absence of neutral currents. However, these would be of the fourth order in the Feynman diagrams of Fig. 1 and second order in the contact interaction form of Eq. (6a). In order to demonstrate the negligibility of the charge-exchange process (for atomic structure), we will use the contact or low-energy form of Eq. (6a), as it lends itself immediately to the well-known results of second-order perturbation theory. The Feynman diagram for the first-order process, i.e., via neutral currents, has the following form (note that the time vector is horizontal):

The Feynman diagram for the charge-exchange second-order process has the following form:

Notice that for protons the end products are the same. The matrix element for the first-order process, Eq. (10b), which we call $\langle M_1 \rangle$, can be represented as follows:

$$\langle M_1 \rangle \equiv \langle e' p' | h_i | e p \rangle \quad (11a)$$

$$\approx \frac{G_F}{\sqrt{2}} (\alpha Z + \beta N), \quad (11b)$$

where we will assume that the constants α and β are of

the order of unity, as in β -allowed β decay. We will do a similar thing for the second-order process, Eq. (10b). The matrix element for this process we call $\langle M_2 \rangle$. Because it is second order, its appearance is more complicated:

$$\langle M_2 \rangle = \sum_c \frac{\langle \alpha' p | h_{ex} | \nu_c n_c \rangle \langle \nu_c n_c | h_{ex} | \alpha p \rangle}{\Delta E_c} \quad (12a)$$

$$\approx \frac{G_\beta^2}{2} \langle \alpha' Z | \left\langle \frac{cp_{nc} W_{nc}}{\Delta E_c (\hbar c)^3} \right\rangle \rangle \quad (12b)$$

where for the sake of clarity we have dropped (for the moment) the normalized notation ($\hbar = c = 1$). Also, we have given the relativistic form of our approximation for $\langle M_2 \rangle$ in (12b). The summation over the intermediate states c is represented by the quantity in the brackets $\langle \rangle$. Again we will assume that α' is of the order of unity, as for α and β in $\langle M_1 \rangle$. It is easy to check that the dimensions of Eqs. (11) and (12) are the same, namely ergs-cm³ (in cgs units). It is easy to make a numerical estimate for $\langle M_1 \rangle$. However, we will now return to normalized units, i.e., $\hbar = c = m_N \equiv 1$, where m_N is the nucleon mass. (We ignore the relatively small neutron-proton mass difference.) The result we then get is

$$\langle M_1 \rangle \approx 10^{-5} A, \quad (13)$$

where A is the atomic number ($N + Z$). A numerical estimate of $\langle M_2 \rangle$ is a little more difficult, since we must make some guess of reasonable values for the nucleon momentum (p_{nc}), total energy (W_{nc}), and borrowed

energy (ΔE_c) in the intermediate "neutral" state $\nu_c n_c$. For stable nuclei it is not hard to show that

$$\begin{aligned} cp_n &\approx 300 \text{ MeV} \\ W_n &\approx 1000 \text{ MeV} \\ \Delta E_c &\approx 10 \text{ MeV} \end{aligned} \quad (14)$$

None of these estimates is much better than a factor of two, but they serve our order-of-magnitude purpose adequately. Using Eqs. (12) and (14), we find

$$\langle M_2 \rangle \approx 10^{-10} \times 10^2 Z = 10^{-8} Z. \quad (15)$$

It is obvious that $\langle M_2 \rangle \approx 10^{-3} \langle M_1 \rangle$, which, as we shall see, means that parity-nonconserving effects due to "ordinary" β decay in atomic structure are not practically observable. The "new" component of the β -decay interaction, however, which leads to $\langle M_1 \rangle$, does lead to the possibility of observable effects, and that is the subject of the following sections. We cannot close, however, without noting that at high enough proton and electron energies $\langle M_1 \rangle$ and $\langle M_2 \rangle$ start to become comparable. [This is why we retained the relativistic-form Eq. (12b).] For electrons impinging on stationary protons, this energy region begins in the region of $\gtrsim 1000$ GeV. For the opposite situation — namely, protons impinging on stationary electrons — the necessary energy becomes astronomical because of the relativistic transformation, $\gtrsim 10^9$ GeV. For many obvious reasons, neither type of experiment is practical, which is why neutrino beams (up till now) have been the main source of information concerning the existence and nature of the neutral currents.⁸ However, the use of colliding beams of positrons and electrons (or other combinations) is now being actively studied as a new approach to the problem.¹⁶

II. Parity Violation in Atomic Structure

It is clear from the structure of Eq. (9b) that there are two sources for the total parity-violating potential (V_V^e), the electron-nucleon (V_V^{eN}), and the electron-electron (V_V^{ee}) parts. Thus,

$$V_V^e \equiv V_V^{eN} + V_V^{ee}. \quad (16)$$

Both contributions were studied by Bouchiat and Bouchiat,⁶ and we will review their results here. We consider first V_V^{eN} [see Eq. (9b)]:

$$V_V^{eN} = \frac{G_\beta}{\sqrt{2}} \sum_{p,n} \left[(1-4 \sin^2 \theta_w) \right. \\ \left. \times \left(\vec{J}_e^\lambda \cdot \vec{J}_p^\lambda + (\sigma_p^k - \sigma_n^k) J_e^k \right) - \vec{J}_e^\lambda \cdot \vec{J}_n^\lambda \right], \quad (17a)$$

where the summation is over $\lambda = 1-4$ and $k = 1-3$ as well as all nucleons ($p \equiv$ proton, $n \equiv$ neutron). It is not difficult to demonstrate in the nonrelativistic limit, for the nucleon currents only, that all terms except the $\lambda = 4$ term can be neglected. Also, we will assume (as do Bouchiat and Bouchiat⁶) that terms of the form $\sigma_{p,n}^k \cdot J_e^k$ sum coherently to an effect of order $\approx \frac{1}{A}$ of the total effect, which (except for light nuclei) is negligible. If we now also note that the momentum transfer $\hbar \vec{k}$ to the nucleons must be such that $KR \ll 1$, where R is the nuclear radius, then the operator for each nucleon becomes 1 (the monopole approximation), so that

$$V_V^{eN} = \frac{G_\beta}{\sqrt{2}} [Z(1-4 \sin^2 \theta_w) + N] \vec{J}_e^4. \quad (17b)$$

If we then further assume a nonrelativistic limit for the electronic wave functions, which is not tenable in the vicinity of the nucleus, we get the final and rather simple result,

$$V_V^{eN} = \frac{G_\beta}{4\sqrt{2}} [Z(1-4 \sin^2 \theta_w) - N] \\ \times \frac{1}{m_e} [(\vec{\sigma} \cdot \vec{p}) \delta(\vec{r}) + \delta(\vec{r})(\vec{\sigma} \cdot \vec{p})]. \quad (17c)$$

where all operators ($\vec{\sigma}$, $\vec{p}$, $\vec{r}$) refer to the electron whose mass we denote by m_e ($\hbar = c = 1$). The nonrelativistic approximation for the electronic matrix element ($\psi_e' \cdot J_e^4 \psi_e$) is a poor one, and the necessary corrections for it are derived and given in Ref. 6.

The electronic contribution to V_V^e is given by

$$V_V^{ee} = -\frac{G_\beta}{\sqrt{2}} (1-4 \sin^2 \theta_w) \sum_e' \vec{J}_e^\lambda \cdot \vec{J}_e^\lambda, \quad (18a)$$

where the $\sum'$ means summation over all electrons except one, in the usual way one derives an effective one-body potential from a two-body potential. If we make the same nonrelativistic approximation for the electrons as we did for the nucleons, we attain the result

$$V_V^{ee} = \frac{G_\beta}{\sqrt{2}} (1-4 \sin^2 \theta_w) [\vec{\sigma} \cdot \vec{p} \rho_e(\vec{r}) \\ + \rho_e(\vec{r}) \vec{\sigma} \cdot \vec{p}], \quad (18b)$$

where $\rho_e(\vec{r})$ is the electronic density for the $Z-1$ "inactive" electrons. [The argument leading to Eq. (18b) is somewhat more complex than for V_V^{eN} because of the necessity of considering the exchange contribution; see Ref. 6.]

We now have the two contributions to the total single-electron parity-violating potential (due to neutral currents) in Eqs. (17c) and (18b). Bouchiat and Bouchiat⁶ then argue that only the mixing of $s1/2$ and $p1/2$ states need be considered, for at least two rather obvious reasons:

(1) The fact that the $\delta(\vec{r})$ in V_V^{eN} will effectively mix *only* such states. All other states (such as $p3/2$ and $d3/2$) tend to vanish more and more rapidly as $Z \rightarrow 0$.

(2) The factor $[Z(1-4 \sin^2 \theta_w) - N]$ in V_V^{eN} , as contrasted to the corresponding factor $(1-4 \sin^2 \theta_w)$ in V_V^{ee} , will give a much heavier weight to the former for $Z \gg 1$. This, of course, is a guess which requires more detailed study to verify (see below).

In any case, since the mixing of $s1/2$ and $p1/2$ states via V_V^{eN} turns out to be proportional to Z^3 , or even a

slightly higher power of Z than 3 (see below), the most practical experiments will in fact be limited to rather high Z , cesium ($Z = 55$) being the best and (presently) most accessible example. Thus if point (2) bears up under further examination, as we shall see it apparently does, we will have $V_V^c \approx V_V^N$ for $Z \gg 1$.

To make the final judgment concerning the negligibility of V_V^c , we must evaluate the mixing matrix elements $\langle M_N \rangle$ and $\langle M_e \rangle$.¹⁷

$$\langle M_N \rangle \equiv \langle ns1/2 | V_V^N | n'p1/2 \rangle, \quad (19a)$$

$$\langle M_e \rangle \equiv \langle ns1/2 | V_V^c | n'p1/2 \rangle. \quad (19b)$$

This begs the question as to whether the explicit evaluation of $\langle M_N \rangle$ and $\langle M_e \rangle$ for other types of states (such as $p3/2$ and $d3/2$) would also lead to the result $\langle M_e \rangle \ll \langle M_N \rangle$ for $Z \gg 1$. There is no guarantee that this would be the case, but for the reasons given above the absolute size of the mixing is the final argument; it is really large enough for practical detection only for $j = 1/2$ states.

It is very easy to achieve a rather simple expression for $\langle M_N \rangle$ [see Eqs. (19b) and (17c)] using nonrelativistic wave functions for the electron wave functions:

$$\begin{aligned} \langle M_N \rangle &= \frac{-3i}{16\pi m_e} \frac{G\beta}{\sqrt{2}} \{ (4 \sin^2 \theta_w - 1) Z + N \} \\ &\times (R_{n0} R'_{n'1})_{r=0}, \end{aligned} \quad (20a)$$

where R_{n0} and $R'_{n'1}$ are electronic wave functions, and $R' \equiv dR/dr$. Note that $\langle M_N \rangle$ is purely imaginary, a result of the time-reversal-invariant nature of V_V^c . It is now necessary to evaluate R_{n0} and $R_{n'1}$. However, the scaling of the electronic Hamiltonian already tells us that $R_{n'1}$ must have one factor of Z more than R_{n0} , and for $\rho \equiv r/(a_0/Z) \ll 1$ we know that $R_{n0} \approx$ constant and $R_{n'1} \equiv$ constant $\times \rho$. Thus, $\langle M_N \rangle$ is (roughly speaking) proportional to at least the second power of Z , provided that the normalization of R_{n0} and $R_{n'1}$ does not produce some unpleasant surprise. In fact, Bouchiat and Bouchiat⁶ generalize a procedure originally introduced by Fermi and Segre¹⁸ and Foldy¹⁹ to derive the following formula for $\rho \ll 1$:

$$R_{ne}(\rho) \approx \frac{2^\ell Z^2 \ell!^{1/2}}{n^{3/2} (2\ell+1)!} \rho^\ell. \quad (20b)$$

Eq. (20b) is an approximate rendering of a more exact result incorporating higher-order corrections,⁶ however, it is sufficient for our purposes since it leads immediately to a numerical result for $\langle M_N \rangle$, i.e., Eq. (19a):

$$\langle M_N \rangle \approx \frac{i\alpha^4 G\beta}{4\sqrt{2}\pi} Z^2 \{ (4 \sin^2 \theta_w - 1) Z + N \} / (nn')^{3/2}, \quad (21)$$

where we have used "electronic" units ($\hbar = c = m_e = 1$), and α is the fine-structure constant. Thus $\langle M_N \rangle$ is in units of $m_e c^2$ (0.511 MeV), and it is quite small, since

$$\alpha^4 G\beta m_e c^2 \approx 10^{-7} \text{ eV}.$$

However, the Z^2 dependence will help bring this quickly to more practical levels, as will be seen.

We now proceed to the evaluation of $\langle M_e \rangle$. Using essentially the same procedures as above, Bouchiat and Bouchiat⁶ get the following (nonrelativistic) result:

$$\begin{aligned} \langle M_e \rangle &= \frac{iG\beta\alpha^4}{2\sqrt{2}} (1 - 4 \sin^2 \theta_w) \\ &\times \int_0^\infty \rho_e(x) x^2 \left[R_{n0} \left(\frac{2}{x} + \frac{\partial}{\partial x} \right) R_{n'1} \right. \\ &\quad \left. - \frac{\partial R_{n0}}{\partial x} R_{n'1} \right] dx, \end{aligned} \quad (22a)$$

$$x = r/a_0. \quad (22b)$$

They then use the Thomas-Fermi result for $\rho_e(x)$, which grossly overestimates $\rho_e(x)$ in the vicinity of $x \lesssim Zr/a_0$, in the integral in Eq. (22a). They find that contributions in regions 1 and 2 are comparable, where these regions are defined as follows:

$$\text{Region 1: } x < \lambda_0 \approx Zr_0/a_0 \approx 1,$$

$$\text{Region 2: } x > \lambda_0. \quad (22c)$$

It is important to note that in region 1 the velocity of the electron is $\gtrsim Z\alpha$. For Cs this is already 0.4. Thus

we are again using nonrelativistic mechanics in a relativistic regime. Nonetheless, these results should be adequate for order-of-magnitude guidance. Accordingly, we may note that the integral in Eq. (22b) may be given approximately in a Z -independent version for $\rho \lesssim 1$ by extracting the following factor using Eq. (20):

$$Z^{9/4} \times Z^{1/2} \times Z^{5/2} \times Z^{-3} = Z^{9/4}, \quad (23)$$

where the $Z^{9/4}$ is from the Thomas-Fermi expression for $\rho_e(r)$, $Z^{1/2}$ from R_{n0} , $Z^{5/2}$ from dR_{n1}/dr , and Z^{-3} from the volume corresponding to $r_0 \approx a_0/Z$. Thus we see $\langle M_e \rangle$ is (roughly) proportional to $Z^{2.25}$, whereas $\langle M_N \rangle$ is proportional to Z^n , where $n > 3$. Bouchiat and Bouchiat⁶ derive an approximate (non-relativistic) formula for $\langle M_e \rangle$ and apply it to the evaluation of the ratio $\langle M_e \rangle / \langle M_N \rangle$, using Eq. (20a). For the specific case of Cs they find that this ratio is $\lesssim 1.6 \times 10^{-2}$. Thus they conclude that for $Z \gtrsim 50$ the electronic contribution to V_e^e should never be more than a few percent, and they ignore it in all further considerations, as we shall do.

The final mixing matrix element between $ns1/2$ and $n'p1/2$ is then Eq. (21), with the necessary relativistic correction factor discussed above. This factor has been calculated⁶ for the outer or least-bound s and p electrons and has been found to be 2.8 for Cs ($Z = 55$) and 9.0 for lead ($Z = 82$), large in both cases and increasing with Z , as might be expected. The final expression for the mixing of opposite parity involves a sum over $n'p1/2$ (for mixing into $s1/2$ states), and $|p1/2\rangle$ is given by

$$|p1/2\rangle \approx i 1.7 \times 10^{-17} \left\{ \sum_n \frac{|e_n e_{n'}|^{3/4}}{\epsilon_n - \epsilon_{n'}} \ln |p1/2\rangle \right\} \\ \times [(4 \sin^2 \theta_w - 1)Z + N]Z^2, \quad (24)$$

where ϵ_n is the electronic energy in Rydberg units ($\alpha^2/2$ for $\hbar = c = m_e = 1$). The numerical coefficient 1.7×10^{-17} seems rather small; however, the important features which make $|p1/2\rangle$ (which is purely imaginary — see above) detectable, at least in prospect, have to do with the proper choice of the target atom and radiative transition to be experimentally studied. This we discuss in the next section. What still requires further comment here is the evaluation of the sum over the mixing states $\ln |p1/2\rangle$. This summation must also include the continuum. For the particular case of Cs ($Z = 55$), $\ln s1/2 = 16s1/2$, with the value for the Weinberg angle θ_w given above [see Eq. (7)], we find, using Eq. (24),

$$|Cs; p1/2/6s1/2\rangle \approx i 0.5 \times 10^{-11} \left\{ \sum_{n'} \frac{|e_n e_{n'}|^{3/4}}{(\epsilon_n - \epsilon_{n'})} \ln |p1/2\rangle \right\} \quad (25)$$

It is interesting to calculate the mixing coefficient for $n' = 6$ and $n' = 7$; these are 1.1 and 0.55 respectively, where we have utilized experimental data and also have ignored the common factor $i 0.5 \times 10^{-11}$ [see Eq. (25)]. For $n' = 2$, this (normalized) coefficient is $\lesssim 0.1$; however, Bouchiat and Bouchiat⁶ state that for $n' \rightarrow \infty$ the coefficient approaches a constant instead of zero, as one would expect from the fact that $|e_n| \rightarrow 0$ as $n' \rightarrow \infty$. This presumably has to do with the way they treat the continuum, which, although it is not quantizable, strictly speaking, can be treated in a quasi-discrete fashion. Details are to be given in a future publication.²⁰

III. Parity-Mixing in Dipole Transitions

In this section we will discuss the effect of parity nonconservation on the two types of dipole transitions, electric (E1) and magnetic (M1). The latter are usually called forbidden in atomic transitions, even when allowed by the usual parity (no change) and spin ($\Delta T = T$) rules. This is because the ratio of the matrix elements $\langle E1 \rangle / \langle M1 \rangle$ has the order of magnitude (r/λ_c) , where r is the dimension of the system and λ_c is the electron Compton wavelength (3.88×10^{-11} cm). For atoms (electrons) this ratio is $\geq 10^3$, while for nuclei (nucleons) it is ≥ 10 , which explains the forbidden characterization for atomic transitions.

We will choose a somewhat different approach from Ref. 6 in treating parity-mixed transitions, in that we will treat only the specific case of single electron states with $j = 1/2$. Also we will consider the specific physical situation characterizing the experimental approach suggested here and in Ref. 6. (Other approaches are possible, of course.) A simple schematic of this approach is shown in Fig. 2. The incident light, both highly monochromatic and circularly polarized, is generated in an appropriate laser source. It is absorbed resonantly by the atoms of the gas or vapor target and then is re-emitted from the resonantly excited state in various modes. One of these modes is selected by the spectrometer-detector, and the count rates for the two possible circularly polarized (CP) incident states, N_+ , are measured. As we shall see, this difference will be proportional to the degree of parity-mixing characterizing the states of interest. The re-emitted radiation will also be circularly polarized in general (regardless of whether we have parity-mixing or not), except at 90° , as shown in Fig. 2. The 90° orientation of the detector is not essential, and the ratio $\rho_\pm$,

$$\rho_\pm \equiv \frac{N_+ - N_-}{N_+ + N_-}, \quad (26)$$

should be independent of angle, with a polarization-insensitive detector. This condition is independent of whether or not there are electric or magnetic fields at the target position. The result for $\rho_\pm$ should also be independent of whether the light is incident from the left or the right (see Fig. 2) in the absence of magnetic

Fig. 2. Schematic diagram for the proposed experiment. The incident light is highly monochromatic and circularly polarized (CP). This is absorbed resonantly by the atoms of the gas target and then re-emitted at a lower frequency. The re-emitted radiation will be (in general) circularly polarized, even in the absence of parity mixing, except at 90° (as shown). The difference in count rate for the two senses of incident polarization, $N_+ - N_-$, is proportional to the degree of parity breakdown.

fields. We will discuss these matters in more detail at the end of this section.

In Eq. (24) we have given the general form for the parity-violating component (pvc) in a $1ns1/2$ state. We give the whole ket-vector now for the initial state (usually the ground state) of an atomic system such as Cs (see Fig. 3), but with the pvc renormalized to unity and multiplied by the overall coefficient δ_5 :

$$|i_5\rangle = |1ns1/2\rangle + i\delta_5(n)|p1/2\rangle, \quad (27a)$$

$$\langle p1/2 | p1/2 \rangle \equiv 1, \quad |\delta_5| \ll 1, \quad (27b)$$

where $\delta_5(n)$ will be a function of n and will be constrained to be real. Single-particle $p1/2$ states are also possible, of course, as we see in Fig. 3. In this case,

$$|i'_p\rangle = |1n'p1/2\rangle + i\delta'_p(n')|s1/2\rangle, \quad (28a)$$

$$\langle s1/2 | s1/2 \rangle \equiv 1, \quad |\delta'_p| \ll 1. \quad (28b)$$

Fig. 3. Level scheme for single-electron states of Cs ($Z = 55$) involved in resonant excitation of $6s1/2$ electron to the $7s1/2$ state, a doubly forbidden (M1) transition. (Hyperfine coupling to the $7/2^+$ nucleus has been omitted.) The incident radiation (2.3 eV) leads to the resonant re-emission of two different cascades of photons, resulting in four photon energies, two in the far infrared (0.84 and 0.91 eV) and two in the near infrared (1.46 and 1.39 eV). The present proposal relates to the latter two emissions, which can be easily separated in commercially available spectrometers (see Fig. 2).

It is not at all necessary that $|\delta_s(n)| = |\delta_p(n)|$, but δ_s and δ_p will both be real numbers (no imaginary component). Equations identical with (27) or (28) can be used to describe the appropriate final states, $|f_s\rangle$ or $|f_p\rangle$. The electric dipole operators for the CP radiation incident along the z-axis are as follows ($\hbar = c = m_e = 1$):

$$(E11) \equiv \sqrt{\alpha} i A_0 e_i Y_{11}(\hat{r}) r, \quad (29a)$$

$$(E1\bar{1}) \equiv -\sqrt{\alpha} i A_0 e_i Y_{11}(\hat{r}) r, \quad (29b)$$

$$\hat{e}_i \equiv \frac{1}{\sqrt{2}} (\hat{e}_x + i\hat{e}_y); \quad e_i \equiv -\frac{1}{\sqrt{2}} (\hat{e}_x - i\hat{e}_y), \quad (29c)$$

where α is the fine-structure constant, the notation $\hat{r}$ refers to unit vectors, $\vec{r} \equiv \hat{r} r$ refers to the electron, and we use the notation (ELM). Also, for convenience, we use the convention $\bar{M} \equiv -M$. The constant A_0 is the field strength (electric or magnetic). We note that the conventional dipole moment operator (D1) is given by

$$(D1) \equiv r Y_{10}(r), \quad (30)$$

in units of $e\hbar/\alpha$. It is easy to prove that for all states of the form of (27a) or (28a) the diagonal moment $\langle x | (D1) | x \rangle$ is identically zero so long as $\delta_x(n)$ is real. This is a reflection of the fact that it is necessary but not sufficient that parity be violated for a nonzero electric dipole moment. Time-reversal violation is also necessary, and this would imply an imaginary part for $\delta_x(n)$. These conditions are not restricted to $j = 1/2$ states only, of course.

The magnetic dipole operators corresponding to Eqs. (29a-c) are as follows:

$$(M11) \equiv -\frac{\sqrt{\alpha}}{2} A_0 \hat{e}_i \sigma_i, \quad (31a)$$

$$(M1\bar{1}) \equiv -\frac{\sqrt{\alpha}}{2} A_0 \hat{e}_i \sigma_i, \quad (31b)$$

$$\sigma_i \equiv \frac{1}{\sqrt{2}} (\sigma_x + i\sigma_y); \quad \sigma_i \equiv -\frac{1}{\sqrt{2}} (\sigma_x - i\sigma_y), \quad (31c)$$

where σ_x and σ_y are the usual Pauli spinors.

$$[\sigma_x, \sigma_y, \sigma_z] = \begin{bmatrix} (01) \\ (10) \end{bmatrix}, \begin{bmatrix} (0-i) \\ (i0) \end{bmatrix}, \begin{bmatrix} (10) \\ (0-1) \end{bmatrix}. \quad (31d)$$

The change in relative signs between Eq. (29b), (31b) as compared to Eqs. (29a), (31a) is a consequence of the opposite behavior of the electric and magnetic dipole operators under reflection.

Although diagonal electric dipole moments *must* vanish, off-diagonal ones need not, as is well known. Thus, for example, the total matrix element for a (predominantly) $\ln_s 1/2$ to $\ln_b p 1/2$ electric-dipole (E1) transition is given by

$$\begin{aligned} \langle M_{ps} \rangle^+ &= \left\{ \langle n_b p 1/2, 1/2 | + i \delta_p(n_b) s | 1/2, 1/2 \rangle \right\} \\ &\times \left\{ \langle E1 \rangle + \langle M1 \rangle \right\} \left\{ \ln_s 1/2, \bar{1}/2 \right\} \\ &+ i \delta_s(n_s) p | 1/2, \bar{1}/2 \rangle \}. \end{aligned} \quad (32a)$$

$$\langle M_{ps} \rangle^- \equiv \langle M_{ps} \rangle^+ \text{ (with all projection quantum numbers reversed)}. \quad (32b)$$

Using Eqs. (29a,b) and (31a,b) we can get a more explicit expression for $\langle M_{ps} \rangle^+$.

$$\begin{aligned} \langle M_{ps} \rangle^+ &= \left\{ \left[\langle n_b | \left(\frac{\sqrt{2}}{\sqrt{3}} Y_{11\beta} - \frac{1}{\sqrt{3}} Y_{10\alpha} \right) \right. \right. \\ &\quad \left. \left. - i \delta_p(n_b) K n_b^+ (Y_{00\alpha}) \right] (i r Y_{11} (t) - (a_1/a)) \right. \\ &\quad \times \left[Y_{00\beta} \ln_s + i \delta_s(n_s) \left(\frac{\sqrt{2}}{\sqrt{3}} Y_{11\alpha} \right) \right. \\ &\quad \left. \left. - \frac{1}{\sqrt{3}} Y_{10\beta} \right] \ln_s^+ \right\} \left\{ \sqrt{a} A_0 \hat{e}_1 \right\}. \end{aligned} \quad (33)$$

where $(\alpha\beta)$ are (up,down) spinors, $\ln_s$ represents radial wave functions for single particles in state n_s , and $\langle n_b^+ \rangle$ represents the normalized sum over all states of the contaminating radial functions (see Eq. (27a)). The same spin-angle dependence is common to all components of the pvc and is extracted as shown. Note that we have also extracted the constant vector $\sqrt{a} A_0 \hat{e}_1$, and we will ignore it henceforth as it has no effect on present considerations. Our final answer is,

$$\begin{aligned} \langle M_{ps} \rangle^+ &= i \langle n_b | r | n_s \rangle \left(\frac{2}{12\pi} \right)^{1/2} \\ &+ \frac{i}{\sqrt{2}} \langle n_b^+ | n_s \rangle \left\{ \delta_p(n_b) - \frac{1}{3} \delta_s(n_s) \right\}, \end{aligned} \quad (34)$$

where, as noted, we have left out the vector $\sqrt{a} A_0 \hat{e}_1$. It is easy to verify that $\langle M_{ps} \rangle^-$ is the same as $\langle M_{ps} \rangle^+$ except for the reversal of the relative sign of the (E1) and (M1) parts. We note in passing that the (M1) part is due entirely to the pvc's in both initial and final states and is, as expected, proportional to a linear combination of δ_s and δ_p .

The cross section for photoabsorption is proportional to $\langle M_{ps} \rangle^+{}^2$ and will depend on the incident polarization. It is not difficult to get the following result for the difference $\langle M_{ps} \rangle^+{}^2 - \langle M_{ps} \rangle^-{}^2$ using Eq. (34) and the stated relation between $\langle M_{ps} \rangle^+$ and $\langle M_{ps} \rangle^-$ [see Eq. (32b)]:

$$\begin{aligned} |\langle M_{ps} \rangle^+|^2 - |\langle M_{ps} \rangle^-|^2 &= \\ \frac{2}{\sqrt{12\pi}} \langle n_b | r | n_s \rangle \langle n_b^+ | n_s \rangle &\left[\delta_p(n_b) - \frac{1}{3} \delta_s(n_s) \right]. \end{aligned} \quad (35)$$

The quantity ρ_z [Eq. (26b)] is the fractional change in cross section in going from one polarization sense to the other. Thus

$$\rho_z \equiv 2 \times \frac{|\langle M_{ps} \rangle^+|^2 - |\langle M_{ps} \rangle^-|^2}{|\langle M_{ps} \rangle^+|^2 + |\langle M_{ps} \rangle^-|^2}. \quad (36a)$$

$$\rho_z = (12\pi) \frac{\langle n_b^+ | n_s \rangle}{\langle n_b | r | n_s \rangle} \left[\delta_p(n_b) - \frac{1}{3} \delta_s(n_s) \right]. \quad (36b)$$

Assuming that $|\langle n_b^+ | n_s \rangle| \approx 1$, $|\delta_p| \approx |\delta_s| \approx 10^{-17}$ (see Sec. II), and $\langle n_b | r | n_s \rangle \approx 10^{-8}$ cm ≈ 300 electronic units, unit length = $\hbar/m_e c$, we find

$$|\rho_z| \approx 10^{-18} \text{ NZ}^2 \quad (36c)$$

It is obvious that this horrendously small number is quite undetectable, even for $Z \approx 92$, in which case

$|P| \approx 10^{-12}$. The answer to this dilemma is (hopefully) in the choice of another type of excitation in combination with the proper targets, as we shall now see.

In the above we chose the situation in which the electric dipole is allowed and the magnetic dipole forbidden by parity conservation. This is the meaning of Eqs. (32a,b) and would be exemplified in Fig. 3 by choosing an exciting radiation of a frequency to match either the $(6p3/2, 6s1/2)$ energy difference or the $(6p1/2, 6s1/2)$ energy difference. The emitted resonance radiation would then be unimodal and exactly the same in energy as the incident radiation. This is *not* what is shown in Fig. 3, where the exciting radiation excites the $7p1/2$ state ($h\nu = 2.30$ eV), and there are two de-exciting modes, each consisting of a cascade of two photons, the first one in the far infrared ($\approx 14,000$ Å or ≈ 0.9 eV) and the second one in the near infrared ($\approx 8,700$ Å or ≈ 1.4 eV). The (M1) excitation process can be treated much the same as the (E1) process was treated earlier, only now we will find that the predominant process, although parity-allowed, will be strongly hindered or "doubly forbidden" by the (M1) selection rule. $\Delta n = 0$, as well as the simple fact of being an (M1). A situation is known in He, where the $2s1/2 \rightarrow 1s1/2$ single-particle (M1) transition is similarly forbidden and leads to the existence of quasi-stable excited states of the form $(2s1/2, 1s1/2)$.²¹ The matrix elements $\langle M_{ss} \rangle^{\pm}$ can be written, in complete analogy to Eqs. (32a,b) and (33),

$$\begin{aligned} \langle M_{ss} \rangle^{\pm} = & \left\{ \langle n_b s 1/2, 1/2 | -i\delta_s(n_b \chi_p 1/2, 1/2) \right. \\ & \times [irY_{11}(\hat{r}) - (s_1/4)] \cdot \{ |n_a s 1/2, 1/2 \rangle \\ & \left. + i\delta_s(n_a) |p 1/2, 1/2 \rangle \} \right\} \end{aligned} \quad (37a)$$

$$\begin{aligned} = & -\frac{1}{2\sqrt{2}} \langle n_b | n_a \rangle \\ & + \left(\frac{2}{12\pi} \right)^{1/2} [\delta_s(n_b \chi n_b' | r | n_a) \\ & - \delta_s(n_a) \langle n_b | r | n_a' \rangle]. \end{aligned} \quad (37b)$$

Comparing Eqs. (34) and (37b) shows a very large-scale difference in the part of $\langle M_{ss} \rangle^{\pm}$ due to the pvc as well as the parity-allowed part. The latter is proportional to $\langle n_b | n_a \rangle$ in Eq. (37b), i.e., the radial

overlap for two different single-particle states ($n_b \neq n_a$), which is zero in the first order (e.g., in a Hartree-Fock set of wave functions). In fact the entire situation is very clearly exhibited by the change in $\rho_{\pm}$, which for Eqs. (37a,b) becomes

$$\rho_{\pm} = \frac{8}{(12\pi)^{1/2}} \frac{[\delta_s(n_b \chi n_b' | r | n_a) - \delta_s(n_a \chi n_b | r | n_a')]}{\langle n_b | n_a \rangle} \quad (38a)$$

Using the same approximation as before, we can again get an order-of-magnitude result.

$$|\rho| \approx \frac{10^{-17} \times 300}{\langle n_b | n_a \rangle} NZ^2. \quad (38b)$$

In Eq. (38b) we can see explicitly the consequence of reversing the roles of parity-allowed and parity-forbidden transitions. The electric dipole matrix element has moved to the numerator, bringing with it an enhancement factor, when all is said and done, of ≥ 10 . The denominator now contains the residue of a truly forbidden (M1) matrix element, $\langle n_b | n_a \rangle$. Bouchiat and Bouchiat⁶ estimate this overlap, which is made nonzero by several effects (hyperfine mixing, configuration mixing, etc.) and find it to be $\approx 10^{-5}$. Thus we can say,

$$|\rho| \approx 3 \times 10^{-10} NZ^2. \quad (39)$$

For Cs, $|\rho| \approx 7 \times 10^{-5}$, which, as we shall see in Sec. IV, should be quite detectable. This change in the value of $|\rho|$ by a factor of 10^8 or 10^9 is, of course, nothing less than spectacular, and is the major reason for the original suggestion by Bouchiat and Bouchiat⁶ to use doubly forbidden (M1) transitions in this type of investigation. Unfortunately it appears (at the moment) that the only practically accessible target is in fact Cs vapor. However, in Sec. IV we will discuss other possibilities, including one mentioned by Bouchiat and Bouchiat⁶ in the high-Z region, which require more advanced but achievable (we believe) techniques.

Earlier in this section we alluded to the effects of electric and magnetic fields at the target site. Such

fields are impossible to avoid completely, and so their effects must be understood. We begin with the case of an arbitrary electrostatic field, $\vec{E}_0$. Ignoring the p.w.'s for the moment, the matrix element $\langle M_{ss}^+ \rangle$ has the following form in the presence of $\vec{E}_0$:

$$\begin{aligned} \langle M_{ss}^+ \rangle^* = & \left\{ (\langle n_b | Y_{00} \alpha + \langle n_b' | \epsilon_1' (Y_{11} \alpha + Y_{1\bar{1}} \alpha) \right. \\ & \left. + \epsilon_0' Y_{10} \alpha) \right\} \langle i r Y_{11} - (\sigma_1/4) | \\ & \times \left\{ Y_{00} \beta \ln_a + [\epsilon_1 (Y_{11} \beta + Y_{1\bar{1}} \beta) \right. \\ & \left. + \epsilon_0 Y_{10} \beta] \ln_a' \right\}. \quad (40) \end{aligned}$$

where $|n_{a,b}\rangle$ and $|n_{a,b}'\rangle$ are normalized radial wave functions of the same type as defined earlier ($|n_{a,b}\rangle$ are single-particle states and $|n_{a,b}'\rangle$ are mixtures), α, β are up-down spinor functions (as before), and $\epsilon_1, \epsilon_0, \epsilon_1', \epsilon_0'$ are the mixing coefficients, which will be proportional to E_0 in this first-order formulation. The first thing about Eq. (40) to be noted is that the matrix element for rY_{11} is zero because of the spin-flipping necessary in order to make the matrix element for σ_1 nonzero. Thus there can be no interference between $\langle E1 \rangle$ and $\langle M1 \rangle$ due to the influence of an electrostatic field, and the difference $(N_+ - N_-)$ is not affected. However, the sum $(N_+ + N_-)$ could be affected because of the incoherent contribution to the $\langle M_{ss}^{\pm} \rangle^2$ of the following $\langle E1 \rangle$ term.

$$|\langle M_{ss}^{\pm} \rangle|^2 = \frac{1}{4\pi} |\epsilon_1 \langle n_b | r | n_a' \rangle + \epsilon_1' \langle n_b' | r | n_a \rangle|^2. \quad (41)$$

Since the $(\epsilon_1, \epsilon_1')$ coefficients are proportional to E_0 (see above), the contribution to $\langle M_{ss}^{\pm} \rangle^2$ is proportional to E_0^2 . Thus it is important to keep the stray electric fields well below a level such that the induced $\langle E1 \rangle$ transition could compete with the doubly forbidden $\langle M1 \rangle$ transition, thus diluting the expected effect. Of course it could be argued that such dilution does not attack the basic validity of the extraction of the numerator of Eq. (38a) from the experimental result $\mu^{\pm}$. This is because the denominator of Eq. (38a) is essentially the reduced matrix element for the $\langle M1 \rangle$ transition, which will have to be taken from experiment anyhow (see

below). Thus some stray electric field is tolerable, although it would becloud the proper extraction of the $\langle M1 \rangle$ matrix element from the experimental results.

It is somewhat surprising to find that electric fields of the order of 1 V/cm can produce enough induced $\langle E1 \rangle$ strength to compete with the doubly forbidden $\langle M1 \rangle$ strength.⁵ The argument to show this is rather simple, however. First we note that the main cause will be the n-forbiddenness rule, which reduces the $\langle M1 \rangle$ matrix element by a factor of 10^{-4} to 10^{-5} (see above). The induced mixing coefficient due to E_0 is approximately $eE_0 \langle n_b | r | n_a \rangle / \Delta E$, which is easily seen to be $\approx 10^{-8} E_0$. The $\langle E1 \rangle$ matrix element in units of Bohr magnetons (μ_B) is then $E_0 \times 10^{-8} \times (2r/\lambda_C)$, where $r \approx 10^{-8}$ cm and λ_C is the electron Compton wavelength. Thus we find,

$$\langle \langle E1 \rangle \rangle_{\text{induced}} \approx 10^{-5} E_0 \mu_B. \quad (42)$$

with E_0 in units of volts per centimeter. Since $10^{-5} \mu_B$ is in the range of the forbidden $\langle M1 \rangle$ matrix elements, we see we are at the danger level at 1 V/cm, as noted earlier (also see Ref. 5).

Besides stray electric fields of external origin, we have to worry about the fluctuating electric fields encountered by an atom as it wanders through the vapor or gas target. For most purposes these fields will be encountered suddenly, which is to say in binary collisions. During such collisions the induced $\langle E1 \rangle$ rates are of course very much augmented. Their influence on the effective dipole transition probability will depend on temperature and pressure, and also on the collective states of the target (see Sec. II). At temperature equilibrium, as in a normal vapor, the effect tends to be relatively high. Bouchiat and Bouchiat^{5,20} estimate that the collision-induced $\langle E1 \rangle$ transition rate is 2.7 times the doubly forbidden $\langle M1 \rangle$ rate in Cs vapor at 1 Torr ($\approx 2 \times 10^{16}$ cm⁻³) and 550°K. Presumably the detector resolution is $\lesssim 0.1$ cm⁻¹ $\approx 10^{-5}$ eV. Bouchiat and Bouchiat have also estimated that the energy width of the collision-induced transition is ≥ 100 times the "Doppler" width ($\Gamma_D \approx 2.2 \times 10^{-6}$ eV). Thus the use of a laser to scan for the $\langle M1 \rangle$ absorption peak will turn the problem of collision-broadening into one of subtracting a quasi-continuous background from the peak of interest. Nonetheless, a signal-to-noise ratio of 1:2.7 at the peak is seldom a trivial situation.

In the method being described here we propose to use an uncollimated atomic beam of neutral Cs atoms transverse to the direction of the laser beam. A detailed discussion of this and other features of the present proposal is given in Sec. IV. Here we will note that the collective state of an atomic beam is very much more ordered than a gas in temperature equilibrium. This will lead to a reduction (in our case) of the Doppler broadening by a factor of 2-3, with a concomitant density reduction of ≈ 100 . Mainly (of course) because of the reduction in density, the signal-to-noise ratio will be $\approx 35:1$. (See Sec. IV for more numerical estimates and specifications.)

The discussion of stray magnetic fields is simpler than that of stray electric fields for two reasons: 1) stray magnetic fields cannot lead to parity admixtures, and 2) they are quite negligible [compared to induced (E1) transitions] in collision broadening. According to Bouchiat and Bouchiat,⁵ a magneto-

static field at the level of 0.1-0.2 G could induce a polarization of $\approx 10^{-5}$ by mixing hyperfine states. Such a field would have to be parallel or antiparallel to the direction of the laser beam, and the mixing would be between states of the same total J_z and different total J . While it is easy to see how hyperfine mixing at the level of $\approx 10^{-5}$ can indeed occur at the level of ≈ 0.1 G, it is difficult to see how this could result in a polarization greater than $\approx 10^{-5}$ at temperatures of the order of 200°K or higher. In fact a strong field, say of the order of 200 G, is more than enough to decouple the nuclear and atomic spins and yet will only produce a polarization of $\mu_B H/kT \approx 10^{-5}$ at room temperature. Thus it seems that strong magnetic fields should not be a serious worry in the detection of polarization due to parity breakdown of the atomic Hamiltonian. In any case, the effect of such fields on ρ_z will change sign when the incident laser direction is reversed, while the effect due to parity nonconservation obviously cannot, as pointed out by Bouchiat and Bouchiat.⁵

IV. Proposed Experiments

In this section we present three proposals for observing the effects of parity nonconservation in atomic transitions. All three are based on Fig. 2, but they differ in the targets or in the characteristics of the lasers and their output beams. Some of the laser features which are needed are just now becoming available. With regard to the target, mention has already been made of using beams of neutral atoms rather than a gaseous vapor in temperature equilibrium, the latter being the type of target proposed by Bouchiat and Bouchiat.⁵ The proposals discussed here will have in common the use of multiple uncollimated atomic beams. For the sake of convenience, we give now numerical values for several pertinent quantities in the particular case of Cs vapor ($Z = 55$) and the $6s1/2 \rightarrow 7s1/2$ transition (see Fig. 3):

$$\begin{aligned} &^{133}\text{Cs: m.p.} = 301.6^\circ\text{K; b.p.} = 951.6^\circ\text{K;} \\ &^{55}\text{Cs: solid density} = 1.9 \text{ g/cm}^3. \end{aligned} \quad (4.3a)$$

$$\begin{aligned} \Gamma_D(700^\circ\text{K}) &\equiv \text{Doppler width} \approx 7.5 \times 10^{-6} \text{ eV} \\ &= 0.02 \text{ cm}^{-1} \approx 600 \text{ MHz}. \end{aligned} \quad (4.3b)$$

$$\begin{aligned} \Delta_H &\equiv \text{hyperfine separations } (E(4^+) - E(3^+)): \\ &\approx 8.3 \times 10^{-6} \text{ eV} = 2000 \text{ MHz} = 0.07 \text{ cm}^{-1} \\ &\quad (7s1/2); \\ &= 3.8 \times 10^{-5} \text{ eV} = 9193 \text{ MHz} = 0.82 \text{ cm}^{-1} \\ &\quad (6s1/2). \end{aligned} \quad (4.3c)$$

$$\Gamma_{M1} (\text{forbidden}) \approx 4 \times 10^{-21} \text{ eV } (7s1/2). \quad (4.3d)$$

$$\Gamma_{E1} (\text{allowed}) \approx 5 \times 10^{-8} \text{ eV } (7s1/2). \quad (4.3e)$$

$$\sigma_{\text{res}}(M1) \approx (\Gamma_{M1}/\Gamma_{\text{las}}) \times 2.5 \times 10^{-10} \text{ cm}^2 \\ (6s1/2 \rightarrow 7s1/2). \quad (4.3f)$$

Typical dye laser resolutions, Γ_{las} (dye):

$$\begin{aligned} 250 \text{ mW: } &4.2 \times 10^{-5} \text{ eV} = 10,000 \text{ MHz} = 0.33 \text{ cm}^{-1}; \\ 100 \text{ mW: } &8.4 \times 10^{-6} \text{ eV} = 2,000 \text{ MHz} = 0.07 \text{ cm}^{-1}; \\ 50 \text{ mW: } &\leq 4 \times 10^{-6} \text{ eV} = 10 \text{ MHz} = 0.0003 \text{ cm}^{-1}. \end{aligned} \quad (4.3g)$$

Equations (4.3d,e) are rough single-particle estimates with the addition of a hindrance factor of 10^4 in the

Fig. 4. Single-electron level structure of Cs ($Z = 55$), exhibiting fine and hyperfine structure, the latter due to the coupling with the nuclear spin of $J = 7/2$. See Fig. 3 for gross single-particle energies and allowed transitions.

(M1) matrix element⁶ in the case of Eq. (43d). The 50- and 100-mW figures in Eq. (43g) are somewhat conjectural at this time. The level energies of Cs are shown in Figs. 3 and 4; in Fig. 4 we exhibit the hyperfine splitting explicitly. Equation (43f) is the peak resonance cross section under the assumption that $\Gamma_{\text{las}} \gg \Gamma_D \gg \Gamma_{M1}$, which we will assume to be the case here (see the 100-mW entry of Eq. 43g), since this corresponds to the most reliable mode (for our purposes) in which to operate typical existing or soon-to-be-available dye lasers. (It is notable that the Doppler width is cancelled out of the equation.)

In Fig. 5 we show a diagram of the type of oven proposed here and the placement of the laser beam with respect to the series of 0.5-mm apertures giving rise to the uncollimated atomic beams which form the more-or-less linear target. The laser beam has been placed 2 mm above the apertures, at which point the

effective density will be $\approx 1 \times 10^{14} \text{ cm}^{-3}$, as compared with $\approx 2 \times 10^{16} \text{ cm}^{-3}$ just inside the apertures. It is possible that the 2-mm distance can be reduced and the operating density (at the apertures) increased in a final design, but for the moment we will abide by these numbers, which are the most realistic we can collect as of now. Because the (single-particle) collision rate goes as the density and also because of the relatively coherent motion of the atomic beam, all collisions are essentially eliminated for the situation depicted in Fig. 5. The coherence of even an uncollimated atomic beam is such that the average relative speed of two atoms is less by a factor of ≈ 2 than that in a gas at thermal equilibrium. To illustrate what this means with regard to the rate of collision-induced transitions, we give the following rough estimate of f_c , the fraction of time spent by any atom in a quasi-diatomic situation:

Fig. 5. Diagram of the proposed molybdenum oven which would generate uncollimated atomic beams through which the laser beam passes. The resonant scattering is detected through wide-angle (≈ 1 sr) glass windows. The temperature inside the molybdenum oven will be in the range of liquid Cs or Rb ($\leq 700^\circ\text{K}$), and the density at the apertures will be $\approx 2 \times 10^{16} \text{ cm}^{-3}$. The heating is accomplished by currents passing through the oven. The target region is formed by a series of 0.5-mm apertures giving rise to an effective target length of about 1 cm.

$$f_c \approx N \cdot \langle v_{\text{rel}} \rangle \cdot \sigma_c \cdot t_c \quad (44)$$

where

$N \equiv$ the atomic density,

$\langle v_{\text{rel}} \rangle \equiv$ the mean two-atom relative velocity,

$\sigma_c \equiv$ the collision cross section of the appropriate type (presumably the hard-sphere scattering cross section),

$t_c \equiv$ the characteristic time for the quasi-diatomic association (something like a transit time).

Assuming that both σ_c and t_c are essentially independent of velocity and density at these low temperatures and pressures, we see that collision-induced transitions should be reduced in the same ratio as the product

$N \langle v_{\text{rel}} \rangle$. This ratio becomes (in our case) ≈ 400 as compared with a gas at a pressure of 1 Torr. Thus, collision-induced processes should cease to be a problem altogether, since the signal-to-noise ratio ought to become 150:1 instead of 1:2.7, as estimated by Bouchiat and Bouchiat⁵ (see above). Unfortunately the power output of presently available dye lasers with linewidths narrow enough to distinguish (for example) the $3^+ \rightarrow 4^+$ transition from the $3^+ \rightarrow 3^+$ transitions (see Fig. 4) is very low (≈ 20 mW). Accordingly we have chosen to use the 100-mW figures (linewidth $\approx 10^{-5}$ eV) of Eq. (43g) as our working example, although the 50-mW figures quoted are very feasible (because of the much reduced linewidth). The 100-mW figures bring our signal-to-noise ratio down to $\approx 35:1$, still a very tolerable regime.

We now proceed to more detailed descriptions of the three experiments we propose. We will give them in what we construe to be the present order of increasing difficulty.

A. THE SIMPLE MULTIPLE-ATOMIC-BEAM EXPERIMENT (Cs, $Z = 55$)

This approach has already been described essentially. The experimental setup is given in block form in Fig. 6. We now give the anticipated specifications for the various components of the experiments.

1. Ar⁺ Laser

This should be a 15-W, continuous-wave (CW) laser with a beam diameter of ≈ 1 mm. The angular divergence must be matched to the dimensions and position of the tunable dye laser which it pumps. Commercial units of this power rating (at minimum) are available.

2. Tunable Dye Laser

For the Cs experiment this is the element which is (at the moment) not commercially available [see Eq. (43g)]. However, units such as the 100- or 50-mW-output types should be commercially available within a few months. Suitable dye lasers can, however, be built by proper conversion of the elements of existing commercial units. The angular divergence of all such units are or will be adequate for all present purposes.

3. Circular Polarizer

This consists of two parts: first a Glan polarizer, which produces a 100% linearly polarized beam, which then goes into the second element, a cylindrical-ring-electrode Pockels cell, wherein the degree and type of polarization is determined by the character of the dc voltage applied. The sense is determined by the sign of the voltage. Of course, suitable optical mountings are required to match the position and orientation of the two elements to each other and to the laser beam. All components are commercially available, and, as a matter of fact, the polarizer is about ready for testing as of this date (November 20, 1974).²²

4. Power Supply for Polarizer

This is a rather special power supply which has to be capable of giving square waves at low frequency (≤ 200 Hz) at various amplitudes from 0 to 8 kV. The

rise time should be $\approx 10^{-8}$ sec. Such units are not available commercially, but such a power supply has been designed and should be completed in the near future.

5. Master Control Circuit

This circuit will house the low-frequency oscillator which will trigger the polarizer power supply at a selectable frequency. It will also gate (alternately) the separate banks of scalars in correspondence with the $\pm$ phase of the incoming laser beam (see Fig. 6). By operating at frequencies of ≈ 200 Hz, the gain fluctuations of all elements of the system, which are characterized by time constants $\gg 1$ sec, will be nullified for all practical purposes. This method of "rapid alternation" has been used in similar experiments in the detection of circular polarization due to parity breakdown in β - γ cascades.²³ In these early experiments the main limitation was not the gain stability but the data-collection rate. In the present experiment it is possible that, for $\rho_{\pm} < 10^{-5}$ (see Sec. IVB, below), gain stability may begin to be a problem. Such circuitry as this is not commercially available but has been designed and constructed at LLL, where it is now operating.

6. Monitoring System

Immediately after the circular polarizer, a partially reflecting mirror is put into the beam to deflect a small portion into the monitor photomultiplier (PM) tube (RCA 8575). All of the circuitry associated with this system will be standard, commercial, fast-pulse (rise time $\approx 2-4$ nsec), and modular (NIM-bin), with the exception of the PM base wiring. The ultrafast circuitry is necessary because of the high rate at which the monitor will run, $\approx 10^6$ counts/sec.

7. Monochromator

The resolution specifications on the monochromator are decided essentially by the structural characteristics of the background on which the peak of interest sits. In our case the background will be mainly dark noise and hence totally achromatic. The peakwidth will be effectively the laser width, or about 10^{-5} eV, using the 100-mW figures from Eq. (43g). This is about the width of the hyperfine separation of the excited $7s1/2$ state [see Eq. (43c)]. Using a high-resolution Fabry-Perot interferometer as a monochromator to single out

defined by the intersection region of the laser beam and the multiple atomic beam. Beams like this have been made with barium, a metal with properties similar to those of Cs and Rb.²⁵ Figure 5 is based on previously used design,^{2,2,5} but other designs are now under study.²⁶ Under the conditions described to this point, the count rate C_d at the monochromator detector due to resonant absorption and re-emission (see Fig. 3) can be estimated as follows:

$$C_d \approx \epsilon n_\gamma L n_a \sigma \quad (45a)$$

where

$$\epsilon \equiv \text{the overall detection efficiency} \\ = 2 \times 10^{-2},$$

$$n_\gamma \equiv \text{the photon flux rate at 100 mW} \\ = 3 \times 10^{17},$$

$$L \equiv \text{the target length} = 1 \text{ cm},$$

$$n_a \equiv \text{the atom density} = 1 \times 10^{14} \text{ cm}^{-3},$$

$$\sigma \equiv \text{the resonant absorption cross section} \\ = 1.2 \times 10^{-25} \text{ cm}^2.$$

Thus

$$C_d \approx (2 \times 10^{-2})(3 \times 10^{17})(1.0)(1 \times 10^{14}) \\ (1.2 \times 10^{-25}) \\ \approx 7 \times 10^4 \text{ counts/sec.} \quad (45b)$$

The final estimate of 7×10^4 counts/sec is a fairly reasonable rate, since we require about 4×10^{10} counts or 1-2 days to measure an effect of $\approx 10^{-4}$ to an accuracy of $\approx 10\%$. By decreasing the distance between the beam apertures and the laser beam by a factor of 2, the data collection time is reduced by a factor of 4 or a little less.

Of course, Eqs. (45a,b) do not include the contribution from collision-induced (E1) transitions, as discussed at the end of Sec. III. Another interfering contribution of considerable potential importance is the fluorescent excitation and subsequent de-excitation by the black-body radiation from the surfaces surrounding the target region.²⁷ In the case of Cs this would correspond to the excitation and subsequent

de-excitation of the $6p1/2$ and $6p3/2$ states of Fig. 3. This particular source of background is extremely sensitive to the temperature of the surroundings as well as the frequency of the excitation probers. This is due to the factor $e^{-h\nu/kT}$ in the black-body radiation formula. For our particular oven design we have included two heat shields between the atomic-beam target region and the oven itself. (These details are not shown in Fig. 5.) This should reduce black-body transition rates so that the signal-to-noise ratio due to this particular source should be $\geq 2:1$. This problem will be far less severe for both Rb (Sec. IVB) and Tl (Sec. IVC). In any case, a good many design parameters, including the optimal source-to-beam distance mentioned earlier, will obviously have to be decided in actual operational trials.

B. STRONG-FIELD SEPARATION OF THE MAGNETIC COMPONENTS WITH ALTERNATING LINEARLY POLARIZED LASER BEAMS (Rb, $Z = 37$).

In this section we present a proposal for utilizing the accidental close equalities between available high-power single-mode laser lines and known forbidden (M1) transitions. This will eliminate the need of an intermediary dye laser and the loss of power and resolution connected therewith. The proposal requires the use of intense magnetic fields transverse to the laser-beam directions in order to use the Zeeman effect to change the quasi-degeneracy into an exact equality with the transition energy, within the laser linewidth (≈ 30 MHz). So far, only one case appears feasible, and that is the $5s1/2$ -to- $6s1/2$ transition in rubidium ($Z = 37$). We show in Fig. 7 the effect of a 79-kG field on the energy difference between the $5s1/2$ (ground state) and $6s1/2$ (excited state) in Rb. (Hyperfine effects are ignored for present purposes.) The unperturbed difference is 4966.82 Å. With a 79-kG field the $5s1/2$ ($m = -1/2$) to $6s1/2$ ($m = +1/2$) transition is shifted by $2\mu_B H = 9.2 \times 10^{-4}$ eV. This then matches a known and rather strong line in the Ar^+ laser emission spectrum, 4965 Å. A typical power rating for this line is ≈ 0.8 W with a linewidth of ≈ 30 MHz or $\approx 1.5 \times 10^{-7}$ eV in the single (longitudinal) mode of operation. Higher power ratings appear to be quite feasible (up to 2X or more).

In Fig. 8 we show a suggested arrangement for achieving an 80-kG field at the target position using

Fig. 7. Diagram of Zeeman splitting of 5s_{1/2} and 6s_{1/2} lines in ⁸⁵Rb in a 79-kG field. Hyperfine coupling and details of the (E1) decay modes have been omitted for the sake of simplicity. The 5s_{1/2} (m = -1/2) to 6s_{1/2} (m = +1/2) transition has been shifted to match the energy of a known and strong monochromatic Ar⁺ laser line at 4965 Å. The ΔE corresponding to the necessary 79-kG field is 9.2×10^{-4} eV (ignoring hyperfine shifts).

superconducting Helmholtz coils, which can be constructed to give on-axis fields of 100 to 150 kG. For the arrangement shown in Fig. 8, fields of about 0.75 times the on-axis field are obtained. Other means may be available for getting the required field intensity, but one important advantage of the arrangement of Fig. 8, should it be practically achievable, is that the same oven can be used for both Cs and Rb because of the close similarity of their physical properties. However, because of the smaller Z value of Rb, the value of $\rho_{\pm}$ [Eq. (26)] is reduced by a factor of 3.5, from 7.0×10^{-5} to 2.0×10^{-5} . This means an increase in the number of required counts (for 10% accuracy) of about a factor of 12, which, however, is partially remedied by the increase in available power by a factor of 8 or more.

As we shall see, several other factors enter the situation, so that the amount of time necessary to accumulate the requisite number of counts for a polarization

measurement with an accuracy of $\approx \pm 10\%$ is about twice that of the Cs experiment, or ≤ 10 days (see Sec. IVA, above).

A fundamental change in the manner in which $\rho_{\pm}$ is measured is necessitated by the fact that we can only excite the $\Delta m = +1$ line, the $\Delta m = -1$ line being unavailable. The change is illustrated in Fig. 8. Instead of using (±) circularly polarized radiation, it is necessary to use linearly polarized radiation, where + and - are distinguished by a rotation of 90°. As it happens, another fortuitous advantage of this scheme is that we can use the same polarizer described in Sec. IVA, except that the bias voltage must be adjusted so that, instead of operating, for example, between -4 and +4 kV, we operate between 0 and 8 kV. The reason why the change of 90° in linear polarization produces a coherent change resulting in a nonzero interference between the (E1) and (M1) amplitudes can be seen by studying the vectors $\vec{E}^{\pm}$ in Fig. 8. For the laser beam

Fig. 8. Suggested arrangement for imposing a 79-kG field transverse to the laser-beam/target region for the Rb experiment (see Sec. IVB). The Helmholtz coils would very likely have to be superconducting coils capable of producing on-axis fields of >100 kG. The two electric vectors $\vec{E}_{\pm}$ represent the two senses of linear polarization, in this case.

incident into the plane of the figure, the (M1) and (E1) amplitudes will have opposite signs in their projections onto the direction perpendicular to the field direction. This is the only direction which gives rise to a $\Delta m = +1$ component for either the (E1) or the (M1) amplitude. For the $\vec{E}^-$ direction the signs of the (E1) and (M1) amplitudes will be the same. Thus, $\rho_{\pm}$ will be given by the same formula as before, Eq. (36b). Unfortunately, we lose a factor of $1/\sqrt{2}$ in the field strength in the direction of interest and another factor of 2 because only one of the two circularly polarized modes making up the useful linearly polarized field can produce $\Delta m = +1$ transitions. (The other one produces $\Delta m = -1$.) The overall loss in transition strength is therefore a factor of 8.

However, we gain back a factor of the ratio of the dye laser width Γ_{las} [see Eqs. (43f, g)] to the effective

Doppler width Γ_{De} , which for an uncollimated beam of Rb will be ≈ 350 MHz. The ratio $\Gamma_{\text{las}}/\Gamma_{\text{De}} \approx 5.32$.

Because of the higher frequency of the exciting radiation in the case of Rb, the cross section for the resonant M1 process is increased by $\sim 8\%$.

Putting all of these factors together, we get the following ratio of the expected count rates for Rb and Cs [see Eq. (45)]:

$$C_d(\text{Rb})/C_d(\text{Cs}) \gtrsim 8 \times 1/8 \times 5.32 \times 1.08 \approx 5.8, \quad (46)$$

where

8 = power increase (commercially available Ar^+ lasers).

1/8 = power loss due to projection on direction of interest and division into two circularly polarized components.

5.32 = linewidth gain factor, $\Gamma_{\text{las}}/\Gamma_{\text{De}}$.

1.08 = factor due to change in wavelength.

We see that to gain the same statistics as for the Cs experiment (10% accuracy) we will need about twice the length of time, or $\lesssim 10$ days. This is, of course, a reasonable length of time, but the estimate is based on several extrapolations from the case of Cs to Rb. For one thing, Eq. (46) does not take account of any change in the relativistic correction factor discussed in Sec. II. In fact, this could reduce the size of $\rho_{\pm}$ by as much as a factor of 2. However, Eq. (46) also does not include any change in the retardation factor for the second forbidden (M1) transition. This factor could well be significantly larger for Rb than for Cs, because, for example, of the reduction of the importance of relativistic contributions or configuration mixing to the forbidden (M1) amplitude.⁶ Another interesting possibility is that the contribution due to hyperfine mixing is reduced to zero (effectively) so long as we select the correct m-projection numbers for the nuclear states which are connected by the absorption of the (M1) radiation. This could be done by small changes in the magnetic field.

In any case, the discussion of these rather technical matters has to be academic at this point. What does seem clear, however, is that, under the assumption of

the achievability of an 80-kG field of reasonable uniformity ($\lesssim 1\%$) and stability ($\lesssim 0.1\%$) over a distance of 1 cm and a diameter of ≈ 1 mm, the Rb experiment looks not only feasible but even simpler (in some regards) than the Cs experiment. Of course *both* experiments should be done if at all possible, since the Z-dependence of $\rho_{\pm}$ will provide invaluable and otherwise inaccessible information regarding the character of the parity-violating potential.

C. FORBIDDEN (M1) ULTRAVIOLET RESONANCE EXCITATION BY CRYSTAL FREQUENCY DOUBLING (TI, $Z = 81$)

There appears to be one fairly clear case of a single-electron ground-state structure very similar to Cs ($Z = 55$) and Rb ($Z = 37$) with higher Z, and that is thallium ($Z = 81$). The ground state in TI has the presumptive configuration of $(6s)^2$ ($7p_{1/2}$), ignoring closed shells, and thus has an atomic spin-parity of 1^+2^- . Both of the stable isotopes have a nuclear spin parity of $1/2^+$ with practically identical magnetic moments ($\mu \approx +1.6 \mu_N$). Thus, except for possible finite nuclear size effects (which should be very small), the hyperfine splittings for the two isotopes, $A = 203$ and $A = 205$, will be indistinguishable, and, at least for our purposes, TI can be considered as mono-isotopic. Also, as pointed out by Bouchiat and Bouchiat,⁶ the spin-orbit coupling does not destroy the orthogonality of different radial wave functions (in this case the $6p_{1/2}$ and $7p_{1/2}$ wave functions) in the nonrelativistic limit. However, the importance (in this case) of configuration mixing could be much greater than for Cs. In any case the size of the (M1) matrix element has been estimated to be comparable to that of Cs, i.e., in the range of 10^2 to $10^4 \mu_B$.⁶ The $\langle E1 \rangle$ amplitude, on the other hand, is expected to be ≈ 10 times bigger because of the multiplication of the factor corresponding to Eq. (21) ($\approx Z^2 N$) and the relativistic factor K_F (see above). As in the other cases, the measurement of the size of the forbidden (M1) amplitude is of fundamental importance in order to properly understand the situation.

In Fig. 9 we exhibit the excitation and decay scheme for the TI experiment. It is clear from this figure that several advantages appear as compared with the Cs experiment, provided that an adequate intensity can be obtained for the incident laser beam.

Fig. 9. Atomic-level diagram for $_{81}\text{TI}$. The single-particle character of the diagram is very much the same as for $_{55}\text{Cs}$ (see Fig. 3), except that in this case the ground state and "pumped" state are in p-orbits and the intermediate state is a $6s_{1/2}$ orbit. Either or both of the two low-lying de-excitation branches (3777 Å , 5532 Å) can be used as the detected scattered radiation.

First, the exciting cross section will be bigger in the ratio $(5395/2927) = 1.8$, assuming the same forbiddenness factor ($\approx 10^{-4}$). Second, the detected radiation will be in the visible or near-ultraviolet and thus

in the region of 30-40% detection efficiency for available PM tubes. Finally, we can use a very-low-resolution type of filter--in fact, the same as proposed for the Cs experiment--and still distinguish the two lower resonantly emitted radiations (5532 Å, 3777 Å). Thus the value of $\rho_{\pm}$ can be measured in two different ways on the same target.

In Fig. 10 we exhibit the elements of the proposed method for getting a flux of 7×10^{13} photons/sec at the wavelength 2927 Å. This was calculated on the basis of a count rate (at the detector) of ≈ 200 counts/sec and a cross section of 1.4×10^{-24} cm², the latter derived from using the 10^{-6} inhibition factor. If the inhibition factor is smaller, then $\rho_{\pm}$ is proportionately bigger, and the time required remains the same, for all practical purposes. (Of course this remark applies to all possible cases.) To measure an effect of the order of 10^{-3} to an accuracy of 10% would require 2×10^8 counts, or 2-4 weeks.

The power required in the final output beam according to the above figures is in the microwatt

range. As shown in Fig. 10, we are assuming that a 15-W Ar⁺ laser pumping a Rhodamine 6G dye laser (with a proper etalon) should be capable of yielding a 300-mW beam with a width of 25 MHz at $\lambda = 5854$ Å. This appears to be well within the realm of present-day laser technology. Such a beam properly focused on a suitable ADP (ammonium dihydrogen phosphate) crystal should yield a doubled-frequency power of ≈ 20 mW, according to the best estimates available. This is about 100 times the required figure. Therefore, there does not appear to be much question of the feasibility of this experiment. It is really a question only of time.

We have not mentioned the target problem for the use of Tl, which is somewhat more difficult than for Cs or Rb, since Tl melts at 577°K and boils at 1467°K. However, it would seem that the only basic change in the design of Fig. 5 that might be required is a change from molybdenum to tungsten as the principal oven material. This and other matters are now under study.

Fig. 10. Block diagram for the Tl experiment. Except for the use of a different dye laser, with the addition of a frequency-doubling crystal, the system is essentially the same as that of Fig. 6, especially after the beam emerges from the uv filter at the desired wavelength (2927 Å). The frequencies at the various points along the beam line are as follows: (a) 5145 and 4880 Å; (b) 5854 Å; (c) 2927 Å.

V. Summary

As of the writing of this section (January 1, 1975), it is already true that some potentially important experimental discoveries have been made since the completion of the previous four sections (before December 1, 1974). The *Scientific American* article by Cline et al.²⁸ summarizes rather nicely the implications of the neutrino and antineutrino beam experiments of the Fermi Laboratory and CERN. Also, we now have first results from PEP (Phase I) experiments at SPEAR, the BNL results, and the ADONE results on the two neutral-particle resonances at 3.1 and 3.7 GeV, both of indeterminably (so far) narrow widths.²⁹ The bearing of these resonances on the unified-field theories is yet to be understood, of course. However, it is certainly clear that the atomic-transition tests proposed by Bouchiat and Bouchiat⁵ and described herein are really quite different from either the neutrino-beam or colliding-beam experiments in the kind of information that can be obtained. In fact it appears that the atomic experiments should be unique.³⁰

As of now at LLL, some of the components for doing the Cs experiment (Sec. IVA) are in the process of either construction or detailed design. There are important and essential exceptions. First and foremost is the 15-W Ar⁺ laser and its associated components. Second is the monochromatic filter-detector (Sec. IVA7). The former is a commercial item which has to be purchased and properly installed. Detailed designs for the latter have yet to be made; however, this is not a difficult problem and should be accomplished shortly. Also it does not involve commercial (capital-equipment) purchases.

A rather weak Ar⁺ laser (<0.6 W) is presently being used in Raman scattering^{31,32} and circular-polarization scattering studies,³³ and it may be possible to measure the (M1) cross sections for the forbidden $s1/2 \rightarrow s1/2$ transitions in Cs and Rb (see Secs. IVA,B) with the present apparatus plus those parts in the various stages of design, construction,

and testing. This would involve the use of mainly equipment borrowed from other experiments and could not begin till about June, 1975, or shortly after. Without the 15-W laser the parity breakdown cannot be done. All the TI experiments will have to wait for the next oven-target, presumably to be made of molybdenum. This applies to both the measurement of (M1) cross sections and the measurement of parity breakdown.

Since the measurement of parity breakdown in the case of the Rb target requires a rather gigantic magnetic field, it appears that our ordering of feasibility of Sec. IV should reverse Secs. IVB and IVC, meaning that the TI experiment might be the easier of the two, practically speaking. At this point this is too difficult to say. There is no question that, given the requisite magnetic field, the Rb experiment is really simpler than either the Cs or the TI experiment in many important ways. The impediment of course is the great expense of building the superconducting coils necessary to produce the requisite magnetic field. It appears that such a device is completely feasible technically and indeed even bigger field-producing magnets have already been built for other applications.³⁴

In any case, at LLL our hope is to have, some time after six months, some results on the (M1) cross sections for both Cs and Rb, provided that the rather minimal extra equipment still lacking becomes available in the relatively near future. We call this Phase I. The (M1) measurements are essential to a clear understanding of the measurement of the parity-breakdown parameter $\rho_{\pm}$ [Eq. (26)], which is Phase II. Since most of the hard work necessary to measure $\rho_{\pm}$ will have been done to get the (M1) cross sections, it is not too optimistic to expect completion of Phase II within six months or so following Phase I, provided that the 15-W Ar⁺ laser and the ancillary facilities and equipment (all standard) become available.

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support, without which the starting phase of this work (see Secs. IV and V) would not have been possible. Finally, we note that Sec. IVC, the projected thallium measurement, is essentially the work of R. C. Harnay.

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