Possibility of Measuring Parity Nonconservation with a Single Trapped Atomic Ion

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(Received 11 January 1993)

A new way of measuring atomic parity nonconservation is proposed that utilizes the remarkable sensitivity inherent in experiments with a single trapped atomic ion. The accuracy may be sufficient to provide a valuable test of electroweak theory. As an illustration, the prospects are analyzed for carrying out such a measurement with Ba^+ by observing the ground state spin rotation induced by an intense laser beam.

PACS numbers: 35.10.Wb, 12.15.Ji, 32.80.Pj, 42.50.Wm

Atomic parity nonconservation (PNC) [1], which arises from exchange of the Z_0 boson between atomic electrons and nucleons, has been measured in optical transitions in a number of elements [2,3]. Detailed atomic structure calculations have permitted the experimental results to be compared with predictions of elementary particle theory. In cesium, the measurement [4] and calculation [5] of PNC are now accurate to better than 2%, and provide an important atomic probe of the electroweak interaction that complements results of high energy experiments. Although the standard model of this interaction accounts very well for all observations thus far, more accurate experiments are needed to determine unambiguously the masses of the predicted top quark and Higgs boson(s), to find out whether there is another generation of quarks and leptons, or to reveal the existence of any new physics near the electroweak energy scale such as technicolor, supersymmetric particles, or additional Z_0 bosons [6]. Further progress with atoms will require improving the measurements and the atomic calculations, or canceling atomic structure uncertainties by comparing very accurate PNC measurements on a string of isotopes of the same element [7–9]. The latter approach would offer the cleanest atomic determination of electroweak physics, limited possibly by nuclear structure uncertainties in the heaviest atoms [9] but otherwise only by experimental accuracy.

In this paper I discuss how the great advances of the past decade in trapping a single atomic ion [10–13] might open up the possibility of very accurate atomic PNC measurements on ions, and I analyze one promising technique for single ions using a PNC spin rotation produced by the "light shift" of the ground state Larmor precession frequency. Likely candidate ions would be single isotopes of alkalilike Ba^+, Sr^+, and Ca^+ that are favorable cases for accurate atomic calculations, and strings of isotopes of these or other ions for canceling atomic structure uncertainties.

A single atomic ion suspended in vacuum in a radio-frequency electric trap can be cooled by laser fields to mK temperatures and confined to an orbit smaller than 0.1 μm at the center of the trap, where it will remain for several hours after the cooling laser fields are turned off. In some ions there are forbidden optical transitions between long-lived electronic states that should be especially sensitive to PNC. These transitions can be driven without Doppler broadening by laser fields focused to very high intensity at the ion, further increasing the sensitivity to PNC. These advantages help to compensate for the absence of the high counting rate provided by dense atomic beams or vapors in current PNC measurements. Using standing wave fields, the spatial and temporal phases at the ion can be adjusted to maximize the PNC transition and discriminate against competing effects. Furthermore, because the time-averaged total electric field acting on the ion—including rf trapping fields and stray DC fields—is necessarily very close to zero, a major potential source of spurious parity mixing [2] is automatically reduced. Finally, trace quantities of an isotope would suffice if a rare species is needed for single-ion measurements on a string of isotopes.

Ba^+ is a good example for illustrating all the major points and is discussed next in some detail, leading to an estimate of the sensitivity to be expected in a PNC measurement. The sensitivity is compared with that of other PNC techniques at the end of the paper.

Barium has nine stable isotopes; five are even-even nuclei, the most favorable types for understanding and calculating nuclear structure corrections to atomic PNC. The electronic energies of the lowest S, P, and D states of Ba^+ are shown in Fig. 1. The cooling laser operates on the $6S_{1/2}$-$6P_{1/2}$ allowed $E1$ absorption line at 493 nm, mistuned slightly to the red of resonance to effect Doppler cooling [14,15]. A "cleanup" laser beam operates on the $6P_{1/2}$-$5D_{3/2}$ transition at 649 nm to avoid

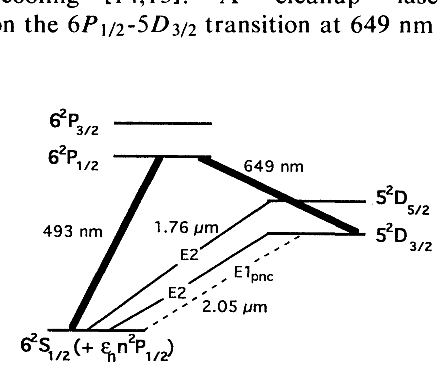

FIG. 1. The Ba^+ energy levels, showing the $E2$ transition at 2.05 μm with the added $E1$ amplitude due to atomic PNC. Also shown are the laser cooling and cleanup transitions at 493 and 649 nm, and the $E2$ shelving transition at 1.76 μm.

losing the ion to the metastable $D_{3/2}$ state. In a typical potential well of 50 eV depth created by rf fields of frequency $\nu_{\rm rf}\cong 25$ MHz, a $\rm Ba^+$ ion oscillating in the trap at a frequency $\nu_{\rm trap}\cong 5$ MHz has been successfully cooled [10,16] to an orbital radius <0.1 μm, with much smaller rf micromotion.

To observe PNC in $\rm Ba^+$ the $E2$-allowed transition between the $6S_{1/2}$ and $5D_{3/2}$ states is of the most interest because of the long lifetime, $\cong 50$ sec, of the excited $5D_{3/2}$ state [17]. The wavelength of this transition, $\lambda=2.05$ μm, is much larger than the size of the cooled ion orbit—a crucial requirement. As depicted in Fig. 1, this transition has a small PNC-induced $E1$ amplitude given by [1,3]

$$\mathcal{E}^{\rm PNC}_{m'm}=\sum_n \frac{\langle 5D_{3/2},m'|e\mathbf{r}|nP_{1/2},m\rangle\langle nP_{1/2},m|H^{\rm PNC}|6S_{1/2},m\rangle}{W_{6S_{1/2}}-W_{nP_{1/2}}}\,, \tag{1}$$

where m and m' denote the magnetic quantum numbers of the two states, W is the electronic binding energy (negative for bound states), and $H^{\rm PNC}$ (a scalar in the electronic variables) describes the short-range electron-nucleon PNC interaction due to Z_0 exchange.

The matrix elements $H^{\rm PNC}_{\bar n n}\equiv\langle\bar n S_{1/2}|H^{\rm PNC}|nP_{1/2}\rangle$ are given by an expression derived by Bouchiat and Bouchiat [1] for alkali atoms, modified slightly here for the case of alkalilike ions of net charge $eZ_{\rm ion}$,

$$H^{\rm PNC}_{\bar n n}\cong i\,1.8\times10^{-17}Z^2Q_WK_r\frac{(W_nW_{\bar n})^{3/4}(a_0)^{1/2}}{(Z_{\rm ion}+1)e}\,, \tag{2}$$

with K_r a relativistic factor $\cong 3$ for Ba ($Z=56$) and $\cong 9$ for Hg ($Z=80$). The binding energies W increase with the net charge, causing the matrix element to increase linearly with $(Z_{\rm ion}+1)$. $Q_W=Z(1-4\sin^2\theta_W)-N$ is the so-called weak nuclear charge, with $\sin^2\theta_W=0.230\pm0.004$ as determined by high energy measurements [6]. The overall factor i appears automatically for a T-even PNC interaction.

The magnitude of $\mathcal{E}^{\rm PNC}_{m'm}$ can be estimated by considering only the contribution of the nearby $6P_{1/2}$ state to the sum in Eq. (1). Using Eq. (2), together with Bates-Damgaard tables [18] for the radial matrix element, the diagonal $(m'=m)$ component is found to be

$$|\mathcal{E}^{\rm PNC}_{mm}|\cong 1\times10^{-11}ea_0\,. \tag{3}$$

A calculation of the complete sum in Eq. (1) using Hartree-Fock wave functions will be published elsewhere [19].

Observable PNC effects can appear through interference of $\mathcal{E}^{\rm PNC}_{m'm}$ with the electric quadrupole amplitude between the same states, the latter defined by

$$(\mathcal{E}^{\rm quad}_{m'm})_{ij}\equiv\left\langle 5D_{3/2},m'\left|\frac{e}{6}(3x_ix_j-r^2\delta_{ij})\right|6S_{1/2},m\right\rangle.$$

In the presence of a light wave having an electric field given by

$$\mathbf{E}(\mathbf{r},t)=\tfrac12[\mathbf{E}(\mathbf{r})e^{-i\omega t}+\text{c.c.}],$$

the PNC and $E2$ couplings are described by the interaction matrix elements

$$\Omega^{\rm PNC}_{m'm}=-\frac{1}{2\hbar}\sum_i(\mathcal{E}^{\rm PNC}_{m'm})_iE_i(0)\,,$$
$$\Omega^{\rm quad}_{m'm}=-\frac{1}{2\hbar}\sum_{i,j}(\mathcal{E}^{\rm quad}_{m'm})_{ij}\left[\frac{\partial E_i(\mathbf{r})}{\partial x_j}\right]_0\,, \tag{4}$$

each evaluated at $r=0$, the location of the ion. Observable quantities are determined by

$$|\Omega_{m'm}|^2=|\Omega^{\rm quad}_{m'm}+\Omega^{\rm PNC}_{m'm}|^2$$
$$\cong|\Omega^{\rm quad}_{m'm}|^2+2{\rm Re}(\Omega^{\rm PNC*}_{m'm}\Omega^{\rm quad}_{m'm})\,, \tag{5}$$

where the small interference term linear in $\Omega^{\rm PNC}_{m'm}$ contains the PNC effects.

If the ion is placed initially in the mth magnetic sublevel of the $6S_{1/2}$ state, the optical field driving the $6S_{1/2}\to 5D_{3/2}$ transition will cause this sublevel to have an energy shift $\hbar\Delta\omega_m$ (called here the *light shift* [20]) as well as a loss rate Γ_m, which together make up the complex level shift $\tilde a_m=\Delta\omega_m-i\Gamma_m/2$. It will be useful to consider optical fields that leave the energy diagonal in m, which requires that the matrix elements in Eq. (4) not connect both $m=+\frac12$ and $-\frac12$ to a common m' sublevel. When in addition the Zeeman splitting among the magnetic sublevels is negligible compared to $|\tilde a_m|$, the field-dependent solution to the set of two-level equations connecting m to the various m' sublevels is [21,22]

$$\tilde a_m=\tfrac12(\omega_0-\omega-i\gamma_{5D}/2)$$
$$\pm\tfrac12[(\omega_0-\omega-i\gamma_{5D}/2)^2+4\Omega_m^2]^{1/2}\,, \tag{6}$$

where $\Omega_m^2=\sum_{m'}|\Omega_{m'm}|^2$, $\hbar\omega_0=W_{5D_{3/2}}-W_{6S_{1/2}}$, and γ_{5D} is the decay rate of the $5D_{3/2}$ state.

For simplicity, it will be assumed henceforth that $\Omega_m\gg|\omega_0-\omega-i\gamma_{5D}/2|$, in which case the real and imaginary parts of Eq. (6) take the forms

$$\Delta\omega_m\to(\omega_0-\omega)/2\pm\Omega_m\,,\quad \Gamma_m\to\gamma_{5D}/2\,, \tag{7}$$

where the light shift shows the well-known sidebands at the "Rabi flopping" frequency. For definiteness, the sideband with the minus sign in Eq. (7) will be retained from now on; it corresponds to the dressed state that evolves adiabatically from the ground state when the fields are turned on *below* resonance ($\omega<\omega_0$) [23].

The light shift will contain a small PNC part, $\Delta\omega_m^{\rm PNC}$, due to the interference term in Eq. (5), and usually a much larger part, $\Delta\omega_m^{\rm quad}$, due to the pure quadrupole term,

$$\Delta\omega_m^{\rm PNC}=2{\rm Re}\sum_{m'}\left[\Omega_{m'm}^{\rm PNC*}\frac{\partial(\Delta\omega_m)}{\partial\Omega_{m'm}^{\rm PNC*}}\right]$$
$$\cong -{\rm Re}\sum_{m'}(\Omega_{m'm}^{\rm PNC*}\Omega_{m'm}^{\rm quad})/\Omega_m^{\rm quad}\,, \qquad (8)$$
$$\Delta\omega_m^{\rm quad}\cong(\omega_0-\omega)/2-\Omega_m^{\rm quad}\,,$$

where $(\Omega_m^{\rm quad})^2\equiv\sum_{m'}|\Omega_{m'm}^{\rm quad}|^2$. When the transition field E is made as large as possible, the energy shift continues to grow in sensitivity to PNC, while the loss rate becomes quite insensitive as Eq. (7) makes evident. Note also that the PNC energy shift, in this limit of large fields, depends linearly on the external field *size* but does not reverse with its *direction*, which is consistent with the well-known theorem [3] that a T-even PNC interaction cannot produce an energy shift linear in the external electric field vector.

It will be convenient to consider two standing wave fields, $\mathbf{E}(\mathbf{r})=\mathbf{E}'(\mathbf{r})+\mathbf{E}''(\mathbf{r})$ (which could be applied using two coherent intersecting beams or simply two components of a single beam, depending upon the fields desired) with E' a large field phased to maximize $\Omega^{\rm PNC}$, and E'' a possibly smaller field to control $\Omega^{\rm quad}$,

$$\mathbf{E}'(\mathbf{r})=\mathbf{E}_0'e^{i\phi'}\cos(\mathbf{k}\cdot\mathbf{r}+\theta')\,,$$
$$\mathbf{E}''(\mathbf{r})=\mathbf{E}_0''e^{i\phi''}\cos(\mathbf{k}\cdot\mathbf{r}+\theta'')\,. \qquad (9)$$

A suitable choice of phases would be $\theta'=0$ and $\theta''=\pm\pi/2$, with $\phi'-\phi''$ adjusted to make $\Omega_{m'm}^{\rm PNC*}\Omega_{m'm}^{\rm quad}$ real.

As an example, if m is specified along the z axis, the fields

$$\mathbf{E}'=\hat{\mathbf{x}}E_0'\cos kz\,,\quad \mathbf{E}''=i\hat{\mathbf{x}}E_0''\sin kz\,,$$
or
$$\mathbf{E}''=i\hat{\mathbf{z}}E_0''\sin kx \qquad (10)$$

will produce $m'-m=\pm1$ quadrupole and PNC dipole transitions. Using such fields in Eq. (4) to obtain the matrix elements for the $6S_{1/2}\rightarrow 5D_{3/2}$ transition needed in Eq. (8) yields

$$\Delta\omega_m^{\rm PNC}=(\pm)1\times10^{-11}\frac{ea_0}{2\hbar}E_0'\,,$$
$$\Delta\omega_m^{\rm quad}=-2\times10^{-4}\frac{ea_0}{2\hbar}E_0''\,, \qquad (11)$$

where $(\pm)$ here indicates that $\Delta\omega_m^{\rm PNC}$ reverses sign with m. The PNC matrix elements are computed as in Eq. (3), using the appropriate Clebsch-Gordon coefficient for each $m'm$, and the quadrupole matrix elements can be obtained from the measured lifetime of the $5D_{3/2}$ state. Note that the fractional PNC light shift, $\Delta\omega_m^{\rm PNC}/\Delta\omega^{\rm quad}\approx(\pm)0.5\times10^{-7}(E_0'/E_0'')$, could in principle be enhanced considerably by making E_0'/E_0'' large.

A convenient way to measure the PNC light shift is to look for a change $\Delta\boldsymbol{\omega}_L^{\rm PNC}$ in the $6S_{1/2}$ Larmor precession frequency, i.e., a change in the Zeeman splitting between the $m=\pm\frac{1}{2}$ sublevels of a given dressed state,

$$(\Delta\omega_L^{\rm PNC})_z\equiv\Delta\omega_{1/2}^{\rm PNC}-\Delta\omega_{-1/2}^{\rm PNC}\,. \qquad (12)$$

A major advantage of using the Larmor frequency for measuring the PNC light shift is that the quadrupole contribution, $(\Delta\omega_L^{\rm quad})_z\equiv\Delta\omega_{1/2}^{\rm quad}-\Delta\omega_{-1/2}^{\rm quad}$, cancels out when $\Delta\omega_m^{\rm quad}$ is independent of m, as is the case for the fields in Eq. (10). Therefore, fluctuations in the optical frequency ω and other sources of variation of $\Delta\omega^{\rm quad}$ do not interfere with the PNC measurement. (There are also other promising variations of the PNC light shift technique, such as using optically resolved Zeeman splittings or Zeeman transitions driven by modulated optical fields [19].)

The sense of PNC Larmor precession reveals the handedness of the interaction, as may be seen by generalizing Eq. (11) to vector form,

$$\Delta\boldsymbol{\omega}_L^{\rm PNC}=\eta\langle2(\mathbf{E}\cdot\boldsymbol{\nabla})\dot{\mathbf{E}}+\mathbf{E}\times(\boldsymbol{\nabla}\times\dot{\mathbf{E}})\rangle_t\,, \qquad (13)$$

where the overdots denote time derivatives and the time average becomes $\langle(\mathbf{E}'\cdot\boldsymbol{\nabla})\dot{\mathbf{E}}''+\boldsymbol{\nabla}(\mathbf{E}'\cdot\dot{\mathbf{E}}'')\rangle_t$ when E' and E'' are optimized for PNC and $E2$ transitions, respectively. If instead $\mathbf{E}$ is a running wave, the time average becomes a vector pointing along the wave propagation vector $\mathbf{k}$. The T-even pseudoscalar η has the approximate magnitude $10^{-15}ea_0\omega/\Omega^{\rm quad}$ in the large field limit considered here.

The change in Larmor frequency due to $\Delta\omega_L^{\rm PNC}$ could be measured in a small static magnetic field directed along z, using a near-resonant rf magnetic field to drive the $\Delta m=\pm1$ Zeeman transition. The shift in resonance would be detected by the method of "shelving" [11,24]. A circularly polarized laser beam tuned to the $6S_{1/2}\rightarrow 5D_{5/2}$ transition at 1.76 μm (see Fig. 1) shelves the ion in the long-lived [11] ($\cong$30 sec) $5D_{5/2}$ state with a probability that depends upon whether a Zeeman transition takes place. Shelving prevents the ions from generating the fluorescence normally observed when the cooling and cleanup laser beams are switched back on; the absence of such fluorescence provides detection of the Zeeman transition with $\cong$50% efficiency. $\Delta\omega_L^{\rm PNC}$ could be distinguished from other shifts through the change in Larmor frequency with changes in the directions and amplitudes that make up the PNC signature in Eq. (13), including the phase angles θ and ϕ in Eq. (9). Competing effects due to imperfections in the optics, etc., have been considered and analyzed [25], and appear to be amenable to adequate controls.

Parity mixing from stray electric fields would be a major concern [2], but the ion should see only a very reduced net dc component of electric field, since the time-averaged force on the ion must vanish inside the trap and the effect of magnetic, optical, and gravitational forces can be made very small. The first-order effects of the trap and stray fields take place at sidebands displaced from the PNC transition by multiples of $\nu_{\rm rf}$ and $\nu_{\rm trap}$, and therefore are not able to produce PNC-like interference. A second-order effect in which these sidebands combine with the Lamb-Dicke sidebands [26] could in principle mimic PNC, but should be small.

To estimate the possible sensitivity of a Ba^+ experiment, assume that an optical field $E_0'=2\times10^4$ V/cm drives the PNC amplitude and, for the moment, a comparable field E_0'' drives the $E2$ amplitude. Such fields could be obtained with 100 mW intracavity beams focused to a 10 μm diameter spot size at the ion. Using these fields in Eq. (11) yields the shift $\Delta\omega^{\text{quad}}/2\pi\cong-2$ MHz, common to each $m=\pm\frac{1}{2}$ level, as well as the PNC Larmor shift defined in Eq. (12),

$$|\Delta\omega_L^{\text{PNC}}|/2\pi\cong0.2\text{ Hz}. \qquad (14)$$

This shift could be increased by using a larger field E', but the off-resonant dipole coupling to the $6P$ and $4F$ levels rapidly becomes important [19]. Without affecting the size of $\Delta\omega_L^{\text{PNC}}$, a much smaller field E'' would reduce $\Delta\omega^{\text{quad}}$, so long as the latter remained much greater than the linewidth of the 2.05 μm laser source.

A single measurement can determine $\Delta\omega_L^{\text{PNC}}=\mathcal{E}^{\text{PNC}}\times E_0'/\hbar$ to within an uncertainty $1/\tau\cong\Gamma\cong\gamma_{5D}/2\cong10^{-2}$ sec^{-1}; simultaneous measurements on N atomic systems over a total observation time $t>\tau$ yield an accuracy given by the approximate expression

$$\frac{\mathcal{E}^{\text{PNC}}}{\delta\mathcal{E}^{\text{PNC}}}\cong\frac{\mathcal{E}^{\text{PNC}}E_0}{\hbar}f\sqrt{N\tau t}\,, \qquad (15)$$

where $\delta\mathcal{E}^{\text{PNC}}$ is the shot noise limited uncertainty and f is an efficiency factor determined by the experimental conditions. The efficiency is limited by the sensitivity to spin polarization and by the fraction of the full exponential decay time τ that can be utilized each measurement cycle. Using the numbers above for Ba^+, with $N=1$ ion and $f\cong0.1$, Eq. (15) yields $\mathcal{E}^{\text{PNC}}$ to a statistical accuracy of 1 part in 500 in a time $t=1$ day.

Equation (15) is of general validity and provides a useful figure of merit for comparing different PNC techniques. In current PNC experiments τ is shortened by radiative or collisional relaxation to $<10^{-7}$ sec, and the optical fields are not as large as the tightly focused field assumed here. Overall, the factor $E_0\sqrt{\tau}$ that appears in Eq. (15) can be larger by $>10^7$ for a single ion compared with conventional atomic beams or vapors, compensating for $N=10^{14}$ atoms or more. PNC experiments with laser cooled or trapped neutral atoms have been proposed [8], and with some elements would offer the interesting possibility of having large N *and* long τ.

A more complete discussion of most points raised here will be published elsewhere [19].

The author wishes to thank H. G. Dehmelt, S. K. Lamoreaux, P. K. Majumder, W. Nagourney, and N. Yu for many useful discussions. This work was supported by NSF Grant No. PHY 9206408.

[1] M. A. Bouchiat and C. C. Bouchiat, Phys. Lett. **46B**, 111 (1974); M. A. Bouchiat and C. C. Bouchiat, J. Phys. (Paris) **35**, 899 (1974).

[2] M. A. Bouchiat and L. Pottier, in *Atomic Physics 9*, edited by R. S. Van Dyck, Jr. and E. N. Fortson (World Scientific, Singapore, 1984), p. 246; E. D. Commins, Phys. Scr. **36**, 468 (1987); P. G. H. Sandars, *ibid.* **36**, 904 (1987); D. N. Stacey, *ibid.* **40**, 15 (1992).

[3] N. Fortson and L. L. Lewis, Phys. Rep. **113**, 289 (1984); E. A. Hinds, in *Atomic Physics 11*, edited by S. Haroche, J. C. Gay, and G. Grynberg (World Scientific, Singapore, 1989), p. 151.

[4] M. C. Noecker, B. P. Masterson, and C. E. Wieman, Phys. Rev. Lett. **61**, 310 (1988).

[5] S. A. Blundell, W. R. Johnson, and J. Sapirstein, Phys. Rev. Lett. **65**, 1411 (1990); V. A. Dzuba, V. V. Flambaum, and O. B. Sushkov, Phys. Lett. A **141**, 147 (1989); A. M. Martensson-Pendrill, J. Phys. (Paris) **46**, 1949 (1985); C. Bouchiat and C. A. Piketty, Europhys. Lett. **2**, 511 (1986).

[6] P. Langacker, M. Luo, and A. K. Mann, Rev. Mod. Phys. **64**, 87 (1992); W. Marciano and J. Rosner, Phys. Rev. Lett. **65**, 2963 (1990).

[7] V. A. Dzuba, V. V. Flambaum, and I. B. Khriplovich, Z. Phys. D **1**, 243 (1986).

[8] C. Monroe, W. Swann, H. Robinson, and C. E. Wieman, Phys. Rev. Lett. **65**, 1571 (1990).

[9] E. N. Fortson, Y. Pang, and L. Wilets, Phys. Rev. Lett. **65**, 2857 (1990); S. J. Pollock, E. N. Fortson, and L. Wilets, Phys. Rev. C **46**, 2587 (1992).

[10] W. Neuhauser, M. Hohenstatt, P. E. Toschek, and H. Dehmelt, Phys. Rev. A **22**, 1137 (1980).

[11] W. Nagourney, J. Sandberg, and H. Dehmelt, Phys. Rev. Lett. **56**, 2797 (1986).

[12] F. Diedrich, J. C. Bergquist, W. M. Itano, and D. J. Wineland, Phys. Rev. Lett. **62**, 403 (1989).

[13] W. Nagourney, N. Yu, and H. Dehmelt, Opt. Commun. **79**, 176 (1990).

[14] D. J. Wineland and Wayne M. Itano, Phys. Rev. A **20**, 1521 (1979).

[15] W. Neuhauser, M. Hohenstatt, P. Toschek, and H. Dehmelt, Phys. Rev. Lett. **41**, 233 (1978).

[16] G. Janik, W. Nagourney, and H. Dehmelt, J. Opt. Soc. Am. B **2**, 1251 (1985).

[17] C. Knab-Bernardini *et al.*, Z. Phys. D **24**, 339 (1992).

[18] D. R. Bates and A. Damgaard, Philos. Trans. R. Soc. London A **242**, 101 (1949).

[19] E. N. Fortson *et al.* (to be published).

[20] W. Happer, Rev. Mod. Phys. **44**, 169 (1972).

[21] S. H. Autler and C. H. Townes, Phys. Rev. **100**, 703 (1955).

[22] C. Cohen-Tannoudji, in *Frontiers in Laser Spectroscopy*, edited by R. Balien, S. Haroch, and S. Liberman (North-Holland, Amsterdam, 1977), p. 3; C. Cohen-Tannoudji and S. Reynaud, J. Phys. B **10**, 345 (1977).

[23] J. L. Carlsten and A. Szöke, Phys. Rev. Lett. **36**, 667 (1976).

[24] H. Dehmelt, Bull. Am. Phys. Soc. **18**, 1521 (1973).

[25] For example, a small degree of circular polarization in $\mathbf{E}'$ added to the ideal plane polarization of Eq. (10) would generate an $E2$ amplitude resembling PNC in some ways. This amplitude can be eliminated by using a suitable optically induced Stark splitting in the $5D$ state. Another example is the presence of $\mathbf{E}'$ or $\mathbf{E}''$ components that produce an energy shift not diagonal in m which can be nulled by measuring all components of $\boldsymbol{\omega}_L$.

[26] W. Neuhauser, M. Hohenstatt, P. Toschek, and H. Dehmelt, Appl. Phys. **17**, 123 (1978).