Linear Stark Shift in Dressed Atoms as a Signal to Measure a Nuclear Anapole Moment with a Cold-Atom Fountain or Interferometer

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We demonstrate theoretically the existence of a linear dc Stark shift of the individual substates of an alkali atom in its ground state, dressed by a circularly polarized laser field. It arises from the electroweak nuclear anapole moment violating P but not T. It is characterized by the pseudoscalar $\xi \mathbf{k} \wedge \mathbf{E} \cdot \mathbf{B}$ involving the photon angular momentum and static electric and magnetic fields. We derive the relevant left-right asymmetry with its complete signature in a field configuration selected for a precision measurement with cold-atom beams. The $3, 3 \rightarrow 4, 3$ Cs transition frequency shift amounts to $7~\mu$Hz for a laser power of ≈ 1 kW at 877 nm, $E = 100$ kV/cm and $B \gtrsim 0.5$ G.

Atomic parity violation has built a bridge between atomic physics and both particle and nuclear physics [1]. The concept of a nuclear anapole moment [2] may be ascribed to toroidal currents [3]. Alternatively, it describes a chiral distribution of the atomic nucleus magnetization, i.e., a nuclear helimagnetism, resulting from the parity violating (PV) nuclear forces exerted inside a stable nucleus [4]. In a stationary atomic state, as a result of T-reversal invariance, only the nuclear spin-dependent PV interaction contributes to observable PV effects, while in forbidden optical transitions the dominant nuclear spin-independent one, associated with the weak nuclear charge Q_W, makes the anapole effect difficult to extract.

Today, there exists only one empirical manifestation of this nuclear anapole moment, though an important nuclear property. It arises from PV measurements performed on the highly forbidden 6S-7S Cs transition, yielding electric dipole transition amplitudes, $E_1^{pv} \simeq 0.8 \times 10^{-11} ea_0$, of a slightly different magnitude (± 2%) for the two hyperfine components $\Delta F = \pm 1$ [5]. The reported result presents significant discrepancies with different theoretical predictions and with other manifestations of the PV nuclear forces [3,4]. On the experimental side, it is notoriously difficult to extract such a small fractional difference from measurements where the control of systematic effects represents a major task. In short, there is a clear need for new independent results both in atomic and nuclear physics. In this latter field, new high precision experiments are already underway, namely, the measurement of the directional gamma-ray asymmetry in the $n + p \rightarrow d + \gamma$ experiment and of the neutron spin rotation in helium [6]. We suggest here a new atomic physics approach using a frequency shift measurement in an atomic fountain or interferometer.

Like previous proposals [7–11], this project is well suited to an atom in its ground state, where the effect of the dominant PV potential due to the weak charge cancels, but it presents important additional advantages. First, instead of relying on the left-right asymmetry in a transition rate, it is based, as in [11], on an energy shift of the atomic states, and its measurement takes advantage of the methods developed for the primary ^{133}Cs standard and atomic interferometers [12,13]. Both domains have now reached an impressive, continuously improving, level of accuracy. Nowadays, beyond their metrological interest, cold-atom fountains and interferometers are becoming undisputable tools, leading to a wide variety of high precision measurements of fundamental interest, such as a lower limit on the time variation of the fine structure constant [14], new tests of Lorentz invariance (LI) [15], possibly tighter constraints on the electron electric dipole moment [16,17], and new tests of the Einstein equivalence principle [14]. Second, this *anapole shift* is both a light shift and a dc linear Stark shift of the $|F, m_F\rangle$ ground state sublevels, violating P but not T. It is quadratic in the parity conserving (PC) atom-radiation field interaction, and linear in the interaction of the nuclear anapole moment with a dc Stark field. Its physical origin then differs totally from the NMR frequency shift predicted in a laser beam [8] and also from the PV light shift involved in a project for a measurement of Q_W on a single trapped Ba^+ ion [18]. In the present project, one must apply both a static electric field and a strong laser field, but its phase is irrelevant. So it is not necessary to localize the atoms inside a cavity unlike the case of previous proposals [9,10].

The nuclear spin-dependent PV interaction is responsible not only for a transition electric dipole but also for a permanent one in an alkali atom ground state, involving the operator $\mathbf{s} \wedge \mathbf{I}$, $\mathbf{s}$ and $\mathbf{I}$ being the spin of the electron and the nucleus, respectively. Thus, in an applied static electric field $\mathbf{E}$, the nuclear anapole moment gives rise to the time-independent P-odd, T-even, interaction $V_{ana} = d_I \mathbf{E} \cdot \mathbf{s} \wedge \vec{I} \equiv -d_I \frac{i}{2} [F^2, \mathbf{s} \cdot \mathbf{E}]$, where $\mathbf{F} = \mathbf{s} + \mathbf{I}$. This identity results from properties of the Pauli matrices, $\vec{\sigma} = 2\mathbf{s}$. The commutator appearing in the last expression shows directly that V_{ana} has matrix elements only between states belong-

ing to different F multiplets. From previously published data [5], the electric dipole moment d_I for the Cs ground state can be estimated to $d_I \simeq 2.36(40) \times 10^{-13} |e| a_0$ [11].

We assume that the atom in its ground state is also perturbed by the electric field of a laser beam of frequency ω detuned from an allowed transition of frequency $\omega_0 = \omega - \delta$. Supposing the beam circularly polarized, we write the associated electric classical field $\mathbf{E}_c(\mathbf{r}, t) = \frac{1}{\sqrt{2}} \mathcal{E}(\mathbf{e}(\xi) \exp(-i\omega t + i\mathbf{k} \cdot \mathbf{r}) + \text{c.c.})$, with $\mathbf{e}(\xi) = \frac{1}{\sqrt{2}} \times (\mathbf{e}_1 + i\xi\mathbf{e}_2)$, $\mathbf{e}_i$ being two orthogonal unit real vectors normal to the wave vector $\mathbf{k}$ of the laser beam and $\xi = \pm 1$ defining the helicity. We write the quantized form of the radiation electric field, assuming that only the laser mode $\mathbf{e}(\xi) \exp(-i\omega t + i\mathbf{k} \cdot \mathbf{r})$ is involved in the radiative transition. The associated annihilation (creation) operator $a(a^\dagger)$ is normalized in such a way that the one photon state obtained by applying $a^\dagger$ upon the vacuum state $|0\rangle$ has its energy $\hbar\omega$ localized inside a volume V. If the laser beam is described by a normalized N photon state $|N\rangle$, its energy density is $\epsilon_0 \mathcal{E}^2 = N\hbar\omega/V$. The combined atom-field system is described in terms of its eigenstates $|i\rangle|N\rangle$, the two kets referring, respectively, to the atom and radiation field states. The coupling of the ground state to the nearly resonant excited state via virtual photon absorption and emission gives the dominant contribution to the light shift, of the $6S_{1/2}$ substates. The Hamiltonian $V_\text{rad} = -\sqrt{\frac{\hbar\omega}{2\epsilon_0 V}}(a\mathbf{e}(\xi) \cdot \mathbf{D} + \text{H.c.})$ describes the atom-field dipolar coupling. Its matrix elements between the ground and excited states $|i\rangle = |nS_{1/2}\rangle|N\rangle$ and $|f\rangle = |n'P_J\rangle|N-1\rangle$ reads $\langle f|V_\text{rad}|i\rangle = -\frac{1}{\sqrt{2}}\mathcal{E}\langle n'P_J|\mathbf{e}(\xi) \cdot \mathbf{D}|nS_{1/2}\rangle$.

We write the lowest-order nonzero modification to the ground state energy involving both perturbations V_ana and V_rad. It is linear in V_ana and quadratic in V_rad:

$$E_\alpha^{(3)} = \langle\alpha|V_\text{ana}RV_\text{rad}RV_\text{rad}|\alpha\rangle + \langle\alpha|V_\text{rad}RV_\text{ana}RV_\text{rad}|\alpha\rangle + \langle\alpha|V_\text{rad}RV_\text{rad}RV_\text{ana}|\alpha\rangle. \quad (1)$$

In this expression, the hyperfine (hf) and Zeeman structures are now included in the description of the atom-field states: $|\alpha\rangle = |nS_{1/2}; F, m_F\rangle|N\rangle$, the resolvent operator $R = (E_n - H)^{-1}$ involves the unperturbed Hamiltonian H and the ground state energies E_n of the combined atom-field system, with hf splitting ΔW. As it will become clear later on, the condition $|\Delta W/\hbar\delta| \ll 1$ has to be satisfied in a realistic experiment. Therefore, the second term in Eq. (1) with an energy denominator $|\hbar\delta|^2$ much larger than those of the other terms $|\Delta W\hbar\delta|$, is neglected

$$E_\alpha^{(3)} = \frac{(F - F')}{\Delta W}\langle\alpha|V_\text{ana}|\alpha'\rangle\langle\alpha'|V_\text{rad}RV_\text{rad}|\alpha\rangle + \text{H.c.} \quad (2)$$

Here the state $|\alpha'\rangle$ differs from $|\alpha\rangle$ just by the change of F, m_F into $F' \neq F$, m_F'. The last factor is the second-order modification to the ground state due to its coupling with the radiation field. First, we limit ourselves to the effect of the laser detuning with respect to the single $P_{1/2}$ excited state.

In this case, V_rad operates between $S_{1/2}$ and $P_{1/2}$, i.e., two electronic states of angular momentum $J = \frac{1}{2}$, and the atom part of the operator involved can be simplified to $V_\text{rad} = \hbar\Omega_1\vec{\sigma} \cdot \hat{\epsilon}$, where $\Omega_1 = d\mathcal{E}/\hbar\sqrt{2}$ is the Rabi frequency, and d the $\Delta m_s = 0$ electric dipole matrix element. The operator R can be replaced by the projection operator $\mathcal{P}_{P_{1/2}} = \sum_m |P_{1/2}, m\rangle\langle P_{1/2}, m|$ times $(\hbar\delta)^{-1}$. A change of F is produced only by the vector part of the tensor operator $V_\text{rad}\mathcal{P}_{P_{1/2}}V_\text{rad}$, which is just $\hbar\Omega_1^2\vec{\sigma} \cdot \xi\mathbf{k}$, if we introduce the photon angular momentum $\xi\mathbf{k} = i(\hat{\epsilon} \wedge \hat{\epsilon}^*)\hbar$. Then, substituting the expressions of V_ana and V_rad into Eq. (2), we obtain the anapole energy shift of a $|F, m_F\rangle$ substate:

$$h\Delta\nu_{F,m_F}^\text{ana} = \frac{d_I\Omega_1^2}{\Delta W\delta} \times \sum_{F',m_F'} (F' - F)\langle F, m_F|\mathbf{E} \cdot \vec{\sigma} \wedge \mathbf{I}|F', m_F'\rangle$$
$$\times \langle F', m_F'|\vec{\sigma} \cdot \xi\mathbf{k}|F, m_F\rangle. \quad (3)$$

Formally, the summation over F' can be extended to all states, although only the states $F' \neq F$ do contribute. By using a closure relation, we obtain the compact expression $\langle F, m_F|\mathbf{E} \cdot \xi\mathbf{k} \wedge \mathbf{I}|F, m_F\rangle$. Taking as the quantization axis the direction defined by the applied magnetic field $\mathbf{B}$, supposed weak ($\mu_B B \ll \Delta W$), we finally obtain

$$h\Delta\nu_{F,m_F}^\text{ana} = 2(F - I)(\xi\hat{k} \wedge \hat{E} \cdot \hat{B})\frac{d_I E\hbar\Omega_1^2}{\Delta W\delta}$$
$$\times \langle F, m_F|I_z|F, m_F\rangle, \quad (4)$$

where $\hat{B}$, $\hat{k}$, and $\hat{E}$ are, respectively, unit vectors parallel to $\mathbf{B}$, $\mathbf{k}$, and $\mathbf{E}$. The I_z matrix elements take the form $\lambda_F m_F$, in Cs $\lambda_F = \frac{7}{8}$ and $\frac{9}{8}$ for $F = 4$ and 3, respectively. We note that shifts of opposite signs are predicted for the two $|F = I \pm \frac{1}{2}, m_F\rangle$ states; hence, both contribute constructively to the frequency shift of the $|\Delta F| = 1$, $\Delta m_F = 0$ transitions which, thus, scales as $2m_F$. If $\mathbf{B}$ defines the quantization axis (see below), then its magnitude has no effect on the anapole shift to second order in V_rad where Eq. (4) is valid.

The presence of the pseudoscalar quantity $\xi\hat{k} \wedge \hat{E} \cdot \hat{B} = \chi$ is the mark of parity violation [19]. If the two fields and $\xi\mathbf{k}$ form a rectangular trihedron, χ takes the value ± 1. Reversals of ξ, $\mathbf{E}$, $\mathbf{B}$, and m_F would be the indispensable operations used to identify the anapole frequency shift. They are all the more helpful since this shift appears superposed on a larger PC light shift associated with the scalar part of the operator $V_\text{rad}\mathcal{P}_{P_{1/2}}V_\text{rad}$. Because of the slight detuning difference between the two F states, the scalar light shift of a $|F, m_F\rangle$ state depends on F to first order in $\Delta W/\hbar\delta$. Thus, the well known shift of the hf splitting is $-\frac{\Delta W\Omega_1^2}{\hbar\delta^2}$, whatever the sign of δ. As for the contribution of the $P_{3/2}$ state, for a detuning $\delta_{3/2}$ comparable to δ, it reinforces that of $P_{1/2}$.

Let us, now, consider the relative contributions to the anapole and the PC light-shifts coming from the vector parts of the operator $V_\text{rad}RV_\text{rad}$ associated with both P_J states. The basic property of the projection operators implying $\mathcal{P}_{P_{3/2}} = \mathbf{1} - \mathcal{P}_{P_{1/2}}$, where $\mathbf{1}$ is the unit operator,

allows us to conclude that they are just opposite. Once weighted by the energy denominators, $\hbar\delta_{1/2} \equiv \hbar\delta$ and $\hbar\delta_{3/2} = -\rho\hbar\delta$, we see that, provided that the detunings be of opposite signs [$\rho > 0$, Fig. 1(b)], both contributions add together and reinforce both the anapole shift [Eqs. (2)–(4)] and the vector PC light shift $\xi\frac{\hbar\Omega_1^2}{\delta}$ by the factor $(1 + \rho)/\rho$. [Note that the parameter ρ is simply related to δ by $(1 + \rho)\delta = \Delta_{FS}$, the fine structure interval, and appears just for writing convenience]. As shown in [20], the light shift induced within the $6S_{1/2}$ sublevels by a circularly polarized beam can be described with the help of a fictitious magnetic field $\mathbf{B}_{ls} = \xi\mathbf{k}\frac{\hbar\Omega_1^2}{|\delta|\mu_B}\frac{1+\rho}{\rho}$. This provides a useful benchmark to compare Zeeman and light shifts. If $\mathbf{k} \cdot \mathbf{B} = 0$, the lowest-order correction to $\Delta\nu_{F,m_F}^{\text{ana}}$ due to the quantization-axis tilt induced by $\mathbf{B}_{ls}$ reads $\frac{\delta\Delta\nu^{\text{ana}}}{\Delta\nu^{\text{ana}}} \approx -\frac{1}{2} \times (\frac{B_{ls}}{B})^2$ for $B \gtrsim B_{ls}$. For $\mathbf{k} \cdot \mathbf{B} = 0$ the ξ-dependence of the hf splitting cancels out.

The anapole shift might be measured in a fountain clock type experiment. A ball of cold atoms in an $|F, m_F\rangle$ state is launched vertically along the applied magnetic field. It passes through a microwave cavity and gets prepared in a coherent superposition of the F, m_F and $F', m_F' = m_F$ states described by the wave function $|\psi\rangle = (|F, m_F\rangle|N\rangle + |F', m_F\rangle|N\rangle)/\sqrt{2}$. After rising a distance L above the cavity center, the atoms reach the interaction region, located close to the apex of their trajectory. The dc field $\mathbf{E}$, the polarized laser beam $\xi\mathbf{k}$, and $\mathbf{B}$ form a rectangular trihedron [Fig. 1(a)]. Because of reversal of the atomic velocity, the very small motional magnetic field cancels. The beam waist w, should be adjusted to the atomic ball radius, ~ 5 mm. The atoms falling back reach the detection region. As usual in a fountain clock the measurement relies on the phase shift of the atomic superposition accumulated during the total time τ_f elapsed since its preparation, $\int \phi(t)dt$. Since the interaction lasts for a limited duration τ_i, we expect the variation of the phase shift for the interaction on and off, $2\pi(\Delta\nu_{F,m_F}^{\text{ana}} - \Delta\nu_{F',m_F}^{\text{ana}})$, to be reduced by the factor $\tau_i/\tau_f \simeq \sqrt{2w/L}$, typically 0.18.

Up to now, Ω_1 has appeared as a free parameter. However, it must be kept in mind that it cannot be chosen arbitrarily large: from the coupling induced between the ground and excited states by the radiation field results an instability of the atomic sample crossing the interaction
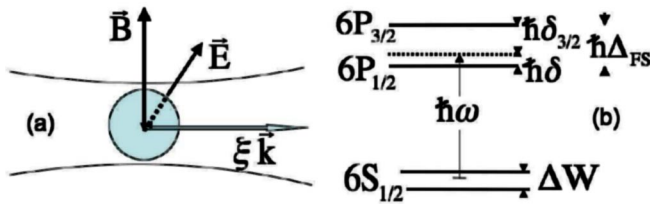

FIG. 1 (color online). (a) Scheme of the interaction region located close to the apex of the cold-atom fountain. (b) Schematics of the relevant Cs atom levels and laser detunings for $\delta_{3/2}/\delta < 0$.

region, hence a signal loss. The decay rate acquired by the 6S atoms is easily computed by perturbation theory in the limit $\Gamma_{PJ} \ll \Omega_1$, $\Gamma_S = \frac{\Omega_1^2}{\delta^2}\Gamma_{P_{1/2}}\kappa$, with $\kappa = 1 + \frac{R_{\text{Rabi}}^2}{\rho^2}\Gamma_{P_{3/2}}/\Gamma_{P_{1/2}}$, where R_{Rabi} is the ratio of the $P_{3/2}$ to $P_{1/2}$ Rabi frequencies. Using Ref. [21], $R_{\text{Rabi}}^2 = 1.98$ and $\kappa = 1 + 2.27/\rho^2$ for Cs. The condition to avoid excessive signal loss is $\Gamma_S\tau_i \lesssim 1$. Assuming this limit just reached, the anapole shift of the $F, m_F \rightarrow F + 1, m_F$ transition angular frequency reads

$$2\pi\Delta\nu_{m_F,m_F}^{\text{ana}} = 2\chi m_F \frac{1 + \rho}{\rho}\frac{d_I E}{\Delta W}\frac{\Omega_1^2}{\delta}\frac{\tau_i}{\tau_f} \tag{5}$$

$$= 2\chi m_F \frac{1 + \rho}{\rho}\frac{d_I E}{\Delta W\tau_f}\frac{\delta}{\kappa\Gamma_{P_{1/2}}} \quad \text{for} \quad \Omega_1^2 = \frac{\delta^2}{\kappa\Gamma_{P_{1/2}}\tau_i}. \tag{6}$$

Equation (6) shows that, due to the constraint imposed on the laser intensity, it is important to choose $\hbar\delta/\Delta W \gg 1$, as announced earlier. Using ρ as the sole free parameter, it is easily found that $\Delta\nu_{m_F,m_F}^{\text{ana}}$ as a function of ρ presents a maximum for $\rho = \sqrt{2.27}$, leading to $\kappa = 2$.

Inserting in Eq. (6), $E = 100$ kV/cm, $d_I E/h = 30$ mHz, $\delta/2\pi = 6.65$ THz, $\Delta W/h = 9.2$ GHz, $\tau_f = 0.5$ s, $\tau_i/\tau_f = 0.18$, $\Gamma_{P_{1/2}}^{-1} = 35$ ns, and $\rho = 1.50$, we obtain for Cs

$$\Delta\nu_{3,3}^{\text{ana}} = 7.6 \ \mu\text{Hz} \quad \text{for} \quad \Omega_1/2\pi = 2.9 \text{ GHz}. \tag{7}$$

leading to $B_{ls} = 0.24$ G. For a beam radius of 5 mm, the laser power should be ≈ 1 kW. This can be achieved inside a Fabry-Perot cavity with a reasonable input power tuned at 877 nm [22]. The condition of adiabatic passage is largely satisfied by the atom-radiation field interaction. One might object that the magnitude chosen for E is large. In fact, still larger fields (450 kV/cm) have been achieved before, but special attention should be paid to the acceleration and defocusing of the atoms entering the field zone [23]. The quadratic differential dc Stark shift [24], $\Delta\nu^{\text{dc}} \simeq 10$ kHz, and the smaller scalar light shift, $\Delta\nu^{ls} = -(1 + \frac{R_{\text{Rabi}}^2}{\rho^2})\frac{\Delta W\Omega_1^2}{h\delta^2}\frac{\tau_i}{\tau_f}$, could be eliminated via the fastest parameter reversals among the set m_F, ξ, $\mathbf{E}$, $\mathbf{B}$.

It is clearly of interest to perform a calibration of the anapole shift allowing for a determination of the basic parameter $d_I E$, free of the uncertainties coming from the laser power and the exact geometry. We suggest using the PC scalar light shift obtained by offsetting sequentially the laser detuning by a fractional amount $\pm\eta$ from one side of its best value to the other. In this case, $\Delta\nu_{\eta-\text{odd}}^{ls} = \Omega_1^2\frac{\tau_i}{\tau_f} \times 0.84\frac{\Delta W}{h\delta^2}\eta$ is the light shift odd in the η reversal for Cs, while $\Delta\nu^{\text{ana}}$, measured at its optimum, is unaffected to first order in η. In the ratio,
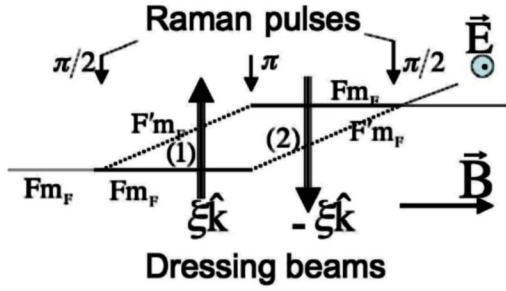
FIG. 2 (color online). Scheme of the interferometer. Both paths cross two dressing beams having opposite photon angular momenta, $\xi\mathbf{k}$.

$$\mathcal{R}(\chi) = \frac{\Delta\nu_{3,3}^{\text{ana}}}{\Delta\nu_{\eta-\text{odd}}^{ls}} = 12\chi\frac{d_I E}{\Delta W}\frac{\hbar\delta}{\Delta W}\eta^{-1}, \qquad (8)$$

which bears the PV signature χ, $\Omega_1^2\tau_i/\tau_f$ is eliminated and the only unknown quantity is just that we need, $d_I E/\Delta W$. The intensity-independent left-right asymmetry $\mathcal{R}(+1) - \mathcal{R}(-1)$ affecting the hf transition frequency is definitely the quantity to be measured. It amounts to $\approx 1.2 \times 10^{-3}$ for $\eta = 0.45 \times 10^{-4}$ (laser frequency excursions of 3 MHz). The precision will depend on the rates of the reversals adjusted to limit the effect of harmful drifts. The most severe of these are likely to be the **B** drifts. Since the $3, 0 \rightarrow 4, 0$ transition is insensitive to the anapole shift, measurements should be performed on a first-order Zeeman shifted transition, as in the case of the LI tests [15]. Relying on frequency comparison with a dual Cs-Rb fountain, already used for a fundamental physics test [14], should reduce the magnetic noise.

Another approach would be to use a cold-atom interferometer. In the case where the wave packets are separated and recombined by using a $\pi/2 - \pi - \pi/2$ sequence of Raman pulses [13], both paths can be made to cross two interaction regions of opposite chiralities (Fig. 2). In the phase difference Φ between the two paths, if **B** is homogeneous the common Zeeman shift cancels, while the anapole shift, odd in ξ reversal, doubles. The same property holds for the calibration signal, if η is applied with opposite signs on the two dressing beams. An estimate gives $\Phi = 4\pi\Delta\nu_{3,3}^{\text{ana}} \approx 10$ μrad for an atom velocity of 1 m/s, $\tau_f = 100$ ms, $\tau_i/\tau_f = 0.2$, Ω_1^2 cut off by 2, for keeping Γ_S as all other parameters the same. Since no absolute measurement is required, this specific phase shift looks accessible in the present state of the art.

In conclusion, we have shown the existence of a linear atomic Stark shift with a chiral character providing a signal for the nuclear anapole moment. We have exhibited a concrete example where several powerful techniques developed recently in atomic physics for high precision frequency determinations appear as highly valuable tools for precise atomic PV measurements. The effect involved presents a certain similarity with the PV NMR frequency shifts expected in enantiomer molecules [25]. The chiral configuration defined by $\xi\mathbf{k}$, **E**, **B**, is playing here the role

of the chiral arrangement of the atoms inside the molecule, but with the advantage of providing numerous reversals. Owing to the absence of any contribution due to the weak nuclear charge, one could thus obtain a rather direct measurement of a PV static property of the nucleus. The method could be extended to other atoms, in particular, to strings of (possibly radioactive) isotopes and provide valuable information on nuclear parity violation.

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