Spins, Electromagnetic Moments, and Isomers of $^{107-120}$Cd

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The neutron-rich isotopes of cadmium up to the $N = 82$ shell closure have been investigated by high-resolution laser spectroscopy. Deep-UV excitation at 214.5 nm and radioactive-beam bunching provided the required experimental sensitivity. Long-lived isomers are observed in $^{127}$Cd and $^{129}$Cd for the first time. One essential feature of the spherical shell model is unambiguously confirmed by a linear increase of the $11/2^-$ quadrupole moments. Remarkably, this mechanism is found to act well beyond the $h_{11/2}$ shell.

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When first proposed the nuclear shell model was largely justified on the basis of magnetic-dipole properties of nuclei. The electric quadrupole moment could have provided an even more stringent test of the model, as it has a very characteristic linear behavior with respect to the number of valence nucleons. However, the scarcity of experimental quadrupole moments at the time did not permit such studies. Nowadays, regardless of experimental challenges, the main difficulty is to predict which nuclei are likely to display this linear signature. The isotopes of cadmium, investigated here, proved to be the most revealing case so far. Furthermore, being in the neighborhood of the “magic” tin, cadmium is of general interest for at least two additional reasons. First, theory relies on nuclei near closed shells for predicting other, more complex systems. Second, our understanding of stellar nucleosynthesis strongly depends on the current knowledge of nuclear properties in the vicinity of the doubly magic tin isotopes. Moreover, specific questions concerning the nuclear structure of the cadmium isotopes require critical evaluation, such as: shell quenching, deformation, sphericity, or whether vibrational nuclei exist at all. Some of these points will be addressed here quite transparently, others require dedicated theoretical work to corroborate our conclusions. In this Letter we report advanced measurements by collinear laser spectroscopy on the very neutron-rich cadmium isotopes. Electromagnetic moments in these complex nuclei are found to behave in an extremely predictable manner. Yet, their description goes beyond conventional interpretation of the nuclear shell model.

The measurements were carried out with the collinear laser spectroscopy setup at ISOLDE-CERN. High-energy protons impinging on a tungsten rod produced low- to medium-energy neutrons inducing fission in a uranium carbide target. Proton-rich spallation products, such as cesium, were largely suppressed in this manner. Further reduction of surface-ionized isobaric contamination was achieved by the use of a quartz transfer line, which allowed the more volatile cadmium to diffuse out of the target while impurities were retained sufficiently long to decay. Cadmium atoms were laser ionized, accelerated to an energy of 30 keV, and mass separated. The ion beam was injected into a gas-filled radio-frequency Paul trap and extracted typically every 100 ms as short bunches with a temporal width of about 5 $\mu$s. The ratio of the above time constants equals the factor of background suppression and therefore results in an increase of the overall sensitivity by the square root of that factor ($\approx 10^2$).

The ion of cadmium was excited in the transition: $5s\ 2^2S_{1/2} \rightarrow 5p\ 2^2P_{3/2}$ at 214.5 nm. Continuous-wave laser light of this wavelength was produced by sequential second-harmonic generation from the output of a titanium-sapphire laser, pumped at 532 nm. The combined fourth-harmonic generation is characterized by a conversion efficiency of up to 2 %. Optimal laser power of about 1 mW was supplied for the measurements. Using the ion for laser excitation increased the overall sensitivity by more than an order of magnitude with respect to the neutral atom. The improvement can be accounted for by the faster transition, the higher quantum efficiency of detection, and the absence of ion-beam neutralization. Such establishment of deep-UV laser beams could potentially provide access to isotopic chains thus far unstudied due to demanding transition wavelengths.
FIG. 1. Example hfs spectra of $^{119}\text{Cd}$ (a), $^{127}\text{Cd}$ (b), and $^{129}\text{Cd}$ (c). Only frequency regions containing hfs components are displayed. The fitted curve incorporates two states on a common background. The lower-spin state is indicated by a semi-transparent fill.

In the conventional manner the atomic hyperfine structure was detected by the ion-beam fluorescence as a function of the laser frequency scanned via the Doppler effect. This method is to a large extent insensitive to contaminant beams. However, care has been taken not to exceed $10^6$ ions accumulated in the Paul trap in order to avoid space-charge effects. This condition was not a limiting factor for the experiment.

An important accomplishment of this work is the discovery of long-lived isomers in $^{127}\text{Cd}$ and $^{129}\text{Cd}$. Representative spectra are displayed in Figs. (b) and (c) where the presence of two nuclear states is clearly identified. It is impossible to determine from the optical measurements alone which of the two is the ground state and what their respective decay modes are. Spins and electromagnetic moments, on the other hand, were determined successfully for both states in each of the isotopes. The presence of such isomers has been suggested in previous studies \[13\,14\].

The experimental results are presented in Tab. I. Some comments on the spin measurements apply here. The hyperfine structure clearly identifies a ground-state spin of $5/2$ for $^{107}\text{Cd}$ and $^{109}\text{Cd}$. Spin $1/2$ is assigned to all ground states from $^{111}\text{Cd}$ to $^{119}\text{Cd}$ due to the reduced number of hfs components, three instead of six. A typical example is $^{119}\text{Cd}$ in Fig. I (a), whose spin adopted in the literature \[16\] is therefore incorrect. The $3/2$ assignments in $^{121-129}\text{Cd}$ are strongly supported by $\chi^2$ analysis of relative hfs intervals and line intensities. Furthermore, the magnetic moments are consistent with an odd-neutron occupation of the $d_3/2$ orbital. The hyperfine structure offers limited sensitivity to high spins. Nevertheless, all $11/2$ assignments are rather firm, since the corresponding electromagnetic moments in Fig. 2 are clearly of $h_{11/2}$ origin.

The $S_{1/2}$ hyperfine parameters $A$ are measured with precision on the level of detectable hyperfine anomaly. Accurate results were deduced with the following procedures. For the observed spins of $1/2$, $5/2$ and $11/2$, there are isotopes in the cadmium chain studied by NMR. The hyperfine anomaly within a set of states with identical spins was neglected and each set was assigned a high-precision value of the corresponding spin as a reference. The resulting magnetic moments are in good agreement with NMR measurements, as evident from Tab. I. For the $3/2$ magnetic moments a hyperfine-anomaly correction was applied with the semi-empirical approach of Moskowitz and Lombardi \[17\]:

$$\frac{A}{A_0} \cdot \frac{I}{I_0} \cdot \frac{\mu_0}{\mu} - 1 = \frac{\alpha}{|\mu_0|} - \frac{\alpha}{|\mu|}. \quad (1)$$

Quadrupole moments were derived from the hyperfine parameters $B$ using the relation: $B = e Q V_{jj}$, where $V_{jj}$ is the electric field gradient at the nucleus and $e$ is the electron charge. Dirac-Hartree-Fock calculations provided the field gradient in the $5p^2 P_{3/2}$ state of the cadmium ion. The finite-difference code GRASP was used to generate the numerical-grid wave functions in conjunction with tools and methodology for hyperfine-structure applications previously described \[20\,21\]. The theoretical error bar was evaluated by applying several methods of orbital generation. Details on the applied computational procedure will be published elsewhere. The obtained electric field gradient is presented in Tab. I along with the quadrupole moments thus determined independent of previous studies. Note that the literature values are about $14\%$ larger in magnitude as they are all referenced to a semi-empirical calculation of the electric field gradient for $^{109}\text{Cd}$ \[22\]. Much of this discrepancy can be accounted for by the Sternheimer shielding, which
with maximum angular-momentum projection, therefore
the magnetic substates $m = \pm j$ are not available for
eucleon pairs. This will produce a quadrupole moment
dependent on the number of nucleons $n$. Since the
number of $j^n$ configurations is $(2j + 1)/2$, Eq. (2)
could explain the alignment of only 6 quadrupole moments
for spin $11/2$. Furthermore, the possibility of configurations
with different seniorities following the same trend can
be excluded. For instance, the matrix element $\langle j^n | \hat{Q} | j'^n \rangle$
for seniority $\nu = 3$, calculated with the aid of tabu-
lated coefficients of fractional parentage [24], is $-8\%$
of the single-particle quadrupole moment $Q_{\text{sp}} = \langle j | \hat{Q} | j \rangle$.
Such values would greatly deviate from the experimen-
tal trend. Clearly, one has to surrender the integer na-
ture of $n$ and interpret it as the actual neutron occupa-
tion. This is possible if one assumes that the population
of neutron pairs ($I = 0$) is shared between the neigh-
boring orbitals: $s_{1/2}$, $d_{3/2}$, $d_{5/2}$, and $h_{11/2}$, suggesting
a kind of degeneracy in terms of total energy per pair.
The odd particle, on the other hand, must always oc-
cupy $h_{11/2}$, as migration to any other orbital in the shell
would result in a change of parity. Finally, assuming no
particle-hole excitations across $N = 82$, one can substi-
tute: $n = 1 + p(A - n_0)$, where $n_0 = 111$, $A = N + Z$, and
$p = 5/9$. The probability $p$ for pair occupation of $h_{11/2}$
is calculated as the capacity of $h_{11/2}$ for neutron pairs
in addition to an odd neutron, divided by the number of
pairs filled between $^{111}\text{Cd}$ and $^{129}\text{Cd}$. It can be easily
verified that with this substitution there is exactly one
$h_{11/2}$ neutron in $^{111}\text{Cd}$ and eleven in $^{129}\text{Cd}$. An exami-
nation of Eq. (4) shows that the quadrupole moments
should cross zero in the middle of the shell, which in the
current description corresponds to $A = 120$. Indeed, the
crossing point was determined at $^{121}\text{Cd}$, very close to that
prediction. In order to account for the small deviation
of one mass unit, the data in Fig. 2 (b) are fitted with an
offset term $Q_{\text{const}}$ representing a constant quadrupole-
moment contribution from correlations with the core.
The resulting fit parameters are: $Q_{\text{sp}} = -667(31)$ mb
and $Q_{\text{const}} = -85(8)$ mb. For comparison, the single-
particle quadrupole moment of $h_{11/2}$ can be estimated
by: $-\langle r^2 \rangle (2j - 1)/(2j + 2) = -269$ mb. Here, under
the assumption of a uniformly charged spherical nucleus,
the mean square radius of the last orbital is approximated
by $5/3$ of the mean square charge radius of $^{111}\text{Cd}$ [30].
The ratio of the two values implies a relatively large ef-
ective charge $e_n = 2.5e$. This result is commented on
below in connection with the magnetic moments. The
line of quadrupole moments crossing zero essentially in
the middle of the $h_{11/2}$ shell indicates a spherical shape
for the $11/2^-$ states. However, one has to acknowledge
the deviation from the straight line at $^{127}\text{Cd}$. It is a small
negative effect occurring between $^{126}\text{Cd}$ and $^{128}\text{Cd}$, for
which abnormal first $2^+$ energies are reported [3]. The
meaning of this observation should be further evaluated
in light of possible shell quenching [3, 6] against sug-

\[ \langle j^n | \hat{Q} | j'^n \rangle = \frac{2j + 1 - 2n}{2j + 1 - 2\nu} \langle j' | \hat{Q} | j \rangle. \] (2)

The linear behavior of the $11/2^-$ quadrupole moments
is the most striking and revealing feature of the cad-
mium nuclei. Moreover, the trend is found to persist
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TABLE I. Spins, hyperfine parameters, and electromagnetic moments derived from this work. Experimental uncertainties (uncorrelated) are quoted in parentheses. Uncertainties on the quadrupole moments due to the electric field gradient (correlated) are enclosed in square brackets. Correction for the hyperfine anomaly is applied to the magnetic moments by using separate NMR references for the states with spin $1/2$, $5/2$, and $11/2$, and by the Moskowitz-Lombardi rule \[15\] for the states with spin $3/2$. High-precision magnetic moments calculated from NMR frequency ratios \[23,20\] relative to the proton \[22\] and for diamagnetism \[25\] are displayed for comparison.

| $Z + N$ | $I$ | $A_{(3p^2_{3/2})}$ (MHz) | $A_{(3p^2_{1/2})}$ (MHz) | $\mu$ ($\mu_N$) | $\mu_{(NMR)}$ ($\mu_N$) | $B_{(3p^2_{3/2})}$ (MHz) | $Q$ (mb) |
|---------|-----|-----------------------------|-----------------------------|-----------------|--------------------------|-----------------------------|---------|
| 107     | 5/2 | $-82.3$ (3)                 | $-3009.8$ (7)               | $-0.6151$ (2)   | $-0.6150554$ (11)        | $401$ (2)                   | $601$ (3) |
| 109     | 5/2 | $-111.4$ (2)                | $-4051.0$ (7)               |                | $-0.8278461$ (15)        | $403$ (1)                   | $604$ (1) |
| 111     | 1/2 | $-398.2$ (5)                | $-14535.0$ (23)             |                | $-0.5948861$ (8)         | $-498$ (3)                  | $-747$ (4) |
| 111     | 1/2 | $-67.2$ (5)                 | $-2456.9$ (5)               | $-1.1052$ (3)   |                          | $-408$ (2)                  | $-612$ (3) |
| 113     | 1/2 | $-418.5$ (6)                | $-15208.0$ (23)             | $-0.6224$ (2)   | $-0.6223009$ (9)         |                          |         |
| 113     | 1/2 | $-66.4$ (2)                 | $-2419.3$ (6)               | $-1.0883$ (3)   |                          | $-408$ (2)                  | $-612$ (3) |
| 115     | 1/2 | $-434.1$ (10)               | $-15840.6$ (30)             | $-0.6483$ (2)   |                          |                          |         |
| 115     | 1/2 | $-63.7$ (2)                 | $-2314.2$ (4)               |                |                          | $-317$ (3)                  | $-476$ (5) |
| 117     | 1/2 | $-499.2$ (11)               | $-18168.5$ (32)             |                |                          |                          |         |
| 117     | 1/2 | $-60.8$ (3)                 | $-2217.5$ (8)               | $-0.9975$ (4)   |                          | $-213$ (4)                  | $-320$ (6) |
| 119     | 1/2 | $-615.5$ (13)               | $-22482.0$ (39)             |                |                          |                          |         |
| 119     | 1/2 | $-59.0$ (2)                 | $-2143.3$ (4)               | $-0.9642$ (3)   |                          | $-90$ (2)                   | $-135$ (3) |
| 121     | 3/2 | $139.7$ (15)                | $5106.2$ (34)               | $0.6269$ (7)    |                          | $-183$ (5)                  | $-274$ (7) |
| 121     | 3/2 | $-62.0$ (3)                 | $-2245.3$ (8)               | $-1.0100$ (4)   |                          | $6$ (4)                    | $9$ (6) |
| 123     | 3/2 | $175.5$ (13)                | $6435.6$ (27)               | $0.7896$ (6)    |                          | $28$ (3)                   | $42$ (5) |
| 123     | 3/2 | $-61.7$ (2)                 | $-2226.3$ (5)               | $-1.0015$ (3)   |                          | $90$ (3)                   | $135$ (4) |
| 125     | 3/2 | $193.5$ (7)                 | $7012.6$ (19)               | $0.8603$ (6)    |                          | $139$ (3)                  | $209$ (4) |
| 125     | 3/2 | $-57.0$ (2)                 | $-2077.9$ (4)               | $-0.9347$ (2)   |                          | $179$ (5)                  | $269$ (7) |
| 127     | 3/2 | $195.3$ (12)                | $7159.6$ (31)               | $0.8783$ (7)    |                          | $159$ (3)                  | $239$ (5) |
| 127     | 3/2 | $-52.6$ (3)                 | $-1934.5$ (5)               | $-0.8702$ (3)   |                          | $228$ (7)                  | $342$ (10) |
| 129     | 3/2 | $187.7$ (23)                | $6912.9$ (48)               | $0.8481$ (8)    |                          | $88$ (5)                   | $132$ (7) |
| 129     | 11/2| $-44.1$ (5)                 | $-1570.2$ (11)              | $-0.7063$ (5)   |                          | $380$ (9)                  | $570$ (13) |

Electric field gradient: $eV_{11/2}/h = 666$ (27) (MHz/b)

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$^a$ Magnetic moment used as a reference for the states with the corresponding spin ($\mu_0$ in Eq. 1, with $\alpha = 0 \mu_N$).

$^b$ $\mu_0$ for the $3/2$ states ($\alpha/\mu_N = 1.7 \%$). The experimental uncertainties of the $3/2$ magnetic moments are convoluted with $5 \times 10^{-4} \mu_N$ representing the standard deviation of the scatter when the hyperfine anomaly is neglected and different reference values are used.

The nuclei of cadmium exhibit yet another simple trend. Their $11/2^-$ magnetic moments, as shown in Fig. 2 (a), increase linearly from $^{111}$Cd to $^{129}$Cd. Four isotopes in the range $^{121-127}$Cd make an exception, which appears to be correlated with the spin change of the second long-lived state. Seemingly, this linear dependence is inconsistent with our description of the quadrupole moments since any odd number of nucleons in a single shell would produce the same magnetic moment as a single nucleon $\frac{1}{2} \otimes \frac{1}{2}$. In this respect one may consider $^{129}$Cd where all neutron orbitals apart from a single $h_{11/2}$ hole are fully occupied with no apparent possibility of “configuration mixing” \[32\]. It is then expected that the $11/2^-$ magnetic moment of $^{129}$Cd should be the most consistent one with the single-particle value, yet it deviates the most. Clearly, an accurate description of the cadmium isotopes should account for the two holes in the $Z = 50$ proton core. First-order core polarization does indeed generate a linear $n$ dependence of the magnetic moment \[34\], though higher-order contributions may be important as well \[35\]. The quadrupole moments, on the other hand, are influenced by this proton-core polarization only through the effective charge, whose large value can now be understood.

In summary, advanced laser spectroscopy provided access to the very exotic odd-mass isotopes of cadmium within the $N = 82$ shell. Long-lived $11/2^-$ states are identified in $^{127}$Cd and $^{129}$Cd for the first time. Remarkably, all quadrupole moments associated with the unique-parity $h_{11/2}$ orbital increase linearly with respect to the number of neutrons, as predicted by the extreme shell-model. Yet, this linear trend is found to extend well beyond the single $h_{11/2}$ shell. Interpretation of both magnetic and quadrupole moments is offered in simple terms and in a common framework.

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[1] M. G. Mayer, Phys. Rev. 78, 16 (1950).
[2] M. G. Mayer and J. H. D. Jensen, Elementary Theory of Nuclear Shell Stricture (John Wiley & Sons, Inc., 1955).
[3] H. Horie and A. Arima, Phys. Rev. 99, 778 (1955).
[4] J. Hakala et al., Phys. Rev. Lett. 109, 032501 (2012).
[5] T. Kautzsch et al., Eur. Phys. J. A 9, 201 (2000).
[6] I. Dilhmann et al., Phys. Rev. Lett. 91, 162503 (2003).
[7] M. Dworschak et al., Phys. Rev. Lett. 100, 072501 (2008).
[8] A. Jungclaus and J. L. Egido, Phys. Scr. T125, 53 (2006).
[9] T. R. Rodríguez, J. L. Egido, and A. Jungclaus, Phys. Lett. B 668, 410 (2008).
[10] P. E. Garrett and J. L. Wood, J. Phys. G 37, 064028 (2010).
[11] U. Köster et al., Nucl. Instrum. Methods Phys. Res., Sect. B 266, 4229 (2008).
[12] E. Mané et al., Eur. Phys. J. A 42, 503 (2009).
[13] N. Hoteling et al., Phys. Rev. C 76, 044324 (2007).
[14] F. Naqvi et al., Phys. Rev. C 82, 034323 (2010).
[15] K.-L. Kratz et al., Eur. Phys. J. A 25, s01, 633 (2005).
[16] D. M. Symochko, E. Browne, and J. K. Tuli, Nucl. Data Sheets 110, 2945 (2009).
[17] P. A. Moskowitz and M. Lombardi, Phys. Lett. B 46, 334 (1973).
[18] I. P. Grant, Relativistic Quantum Theory of Atoms and Molecules: Theory and Computation (Springer, 2007).
[19] P. Jönsson, X. He, C. Froese Fischer, and I. P. Grant, Comput. Phys. Commun. 177, 597 (2007).
[20] J. Bieroń, P. Jönsson, and C. Froese Fischer, Phys. Rev. A 60, 3547 (1999).
[21] J. Bieroń, P. Pyykko, and P. Jönsson, Phys. Rev. A 71, 012502 (2005).
[22] J. Bieroń, C. Froese Fischer, P. Jönsson, and P. Pyykko, J. Phys. B 41, 115002 (2008).
[23] N. S. Laulainen and M. N. McDermott, Phys. Rev. 177, 1615 (1969).
[24] A. de-Shalit and I. Talmi, Nuclear Shell Theory (Academic Press, 1963).
[25] R. L. Chaney and M. N. McDermott, Phys. Lett. A 29, 103 (1969).
[26] P. W. Spence and M. N. McDermott, Phys. Lett. A 42, 273 (1972).
[27] J. Bieroń, P. Pyykko, and P. Jönsson, Phys. Rev. A 71, 012502 (2005).
[28] J. Bieroń, C. Froese Fischer, P. Jönsson, and P. Pyykko, J. Phys. B 41, 115002 (2008).
[29] N. S. Laulainen and M. N. McDermott, Phys. Rev. 177, 1615 (1969).
[30] A. de-Shalit and I. Talmi, Nuclear Shell Theory (Academic Press, 1963).
[31] R. L. Chaney and M. N. McDermott, Phys. Lett. A 29, 103 (1969).
[32] P. W. Spence and M. N. McDermott, Phys. Lett. A 42, 273 (1972).
[33] J. Bieroń, P. Pyykko, and P. Jönsson, Phys. Rev. A 71, 012502 (2005).
[34] J. Bieroń, C. Froese Fischer, P. Jönsson, and P. Pyykko, J. Phys. B 41, 115002 (2008).
[35] N. S. Laulainen and M. N. McDermott, Phys. Rev. 177, 1615 (1969).
[36] A. de-Shalit and I. Talmi, Nuclear Shell Theory (Academic Press, 1963).
[37] R. L. Chaney and M. N. McDermott, Phys. Lett. A 29, 103 (1969).
[38] P. W. Spence and M. N. McDermott, Phys. Lett. A 42, 273 (1972).
[39] J. Bieroń, P. Pyykko, and P. Jönsson, Phys. Rev. A 71, 012502 (2005).
[40] J. Bieroń, C. Froese Fischer, P. Jönsson, and P. Pyykko, J. Phys. B 41, 115002 (2008).
[41] N. S. Laulainen and M. N. McDermott, Phys. Rev. 177, 1615 (1969).