This is what explains the existence of the minimum in the variation of the moment of inertia of the n-phase which will be discussed subsequently.

[5)]The division into two "sub-bands" is equally well-founded also from a purely theoretical point of view. Indeed, one does not expect any considerable differences in principle from even-even nuclei in the present case, while it is desirable to exclude the not entirely clear and nonmacroscopic effect of the "lining up" of the odd nucleon. It is well known that for $K_0=\frac{1}{2}$ a similar effect occurs even in the adiabatic region where it is more easily susceptible of being investigated (cf., for example, Ref. 1).

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Possibility of studying the structure of weak neutral currents in optical transitions in heavy atoms

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Effects of parity nonconservation, caused by the weak interaction of neutral electronic vector and nucleonic axial currents, are examined in optical transitions in thallium, lead, bismuth, and cesium. Degrees of circular polarization of the light are calculated for ordinary and strongly forbidden $M1$ transitions.

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1. Experiments on parity nonconservation effects in heavy atoms (first discussed in[1]) are now being conducted by several research groups.[2–6] From our point of view, the most promising of these experiments are those on the rotation of the plane of polarization of light from heavy-metal vapors, as proposed in[7–9]. They have already given a significant limit on the constant for the weak interaction of the neutral vector current (VC) of nucleons with the axial-vector current (AC) of electrons. Parity nonconservation effects in heavy atoms can also be produced by another type of weak interaction, that between nucleonic AC and electronic VC. The investigation of this interaction would not be unimportant even if there were no interaction of electronic AC with nucleonic VC (which is at present the most popular interpretation of the preliminary experimental results). In fact, the absence of an axial part in the neutral electronic current by no means indicates that it is also absent for nucleons. Moreover, neutrino experiments evidently show that the neutral hadronic current contains both a vector and an axial part.

We have previously pointed out[10,11] the possibility of studying the interaction of the nucleonic AC and the electronic VC via the measurement of the angle of rotation of the plane of polarization of the light at frequencies corresponding to optical transitions between individual hyperfine components of levels in heavy atoms. A complicating feature of the observation of this interaction is that its size is smaller than that of the interaction between the nuclear VC and the electronic AC by roughly a factor Z, the nuclear charge, since only the valence nucleon contributes to the nucleonic AC. However, the accuracies attainable in experiments on the optical activity of heavy-metal vapors are already close to that needed to detect the effect in question. In this paper we make a detailed calculation of the size of the effect. In addition to this, we find the contribution of this interaction to the circular polarization of the radiation emitted in strongly forbidden $M1$ transitions.

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We note that other possible manifestations of the interaction between the nuclear AC and the electronic VC in ions, solids, and heavy atoms have been discussed previously in[12-16]

2. The Hamiltonian for the interaction of the electronic VC with the nucleonic AC is

$$H_W = \frac{G\hbar^3}{c\sqrt{2}} \varkappa_N (\bar{u}_e \gamma_\mu u_e)(\bar{u}_N \gamma_\mu \gamma_5 u_N), \qquad (1)$$

where u_e and u_N are electronic and nucleonic spinors, and $G = 10^{-5}/m_p^2$ is the Fermi constant. The essential point of the experiment in question is to determine the constant $\varkappa_N$, which in general may be different for the proton ($\varkappa_p$) and for the neutron ($\varkappa_n$). The matrix element of the Hamiltonian (1) can be written in the following form[16] (the only nonzero matrix element is that between the $s_{1/2}$ and $p_{1/2}$ states of the electron):

$$\langle s_{1/2}|H_W|p_{1/2}\rangle = i\frac{Gm^2\alpha^2Z^2R}{\pi\sqrt{2}}(\nu_s\nu_{p_{1/2}})^{-3/2}\frac{me^4}{2\hbar^2}\frac{2\gamma+1}{3}\varkappa g_I \cdot 2\mathbf{j}_e\mathbf{I}_n. \qquad (2)$$

Here $\gamma = (1-\alpha^2Z^2)^{1/2}$; ν_s and $\nu_{p1/2}$ are the effective principal quantum numbers of the electron; and R is a relativistic factor ($R_{Tl} = 8.5$, $R_{Pb} = 8.9$, $R_{Bi} = 9.4$, and $R_{Cs} = 2.8$). The factor $\varkappa g_I$, which has appeared in Eq. (2) after the calculation of the matrix element from Eq. (1) with the nuclear wave functions, is taken so that in the shell model of the nucleus, which we shall use for the numerical calculations, $\varkappa$ is equal to the appropriate constant for the valence nucleon ($\varkappa_{Cs} = \varkappa_{Tl} = \varkappa_{Bi} = \varkappa_p$; $\varkappa_{Pb^{207}} = \varkappa_n$) and $g_I = \langle\sigma_e$ (of valence nucleon)$\rangle/I_{A_e}$, and is 2 for Tl, $-2/3$ for Pb^{207}, $-2/11$ for Bi, and $-2/9$ for Cs. It can be expected that this approximation will give an accuracy comparable with that of the predictions of the shell model for the magnetic moments of the nuclei, which is 20 percent for Pb^{207}, 30 percent for Cs and Bi, and 70 percent for Tl.[1)]

In the $M1$ transitions we are considering, parity nonconservation manifests itself in the appearance of a small admixture of the amplitude for the $E1$ transition. This makes the radiation have a circular polarization

$$P = -2\,\mathrm{Im}\frac{\langle J'IF'\|E1\|JIF\rangle}{\langle J'IF'\|M1\|JIF\rangle}. \qquad (3)$$

Here **I** is the angular momentum of the nucleus, J and J' are the initial and final angular momenta of the electrons, and $\mathbf{F} = \mathbf{I} + \mathbf{J}$, $\mathbf{F}' = \mathbf{I} + \mathbf{J}'$ are the initial and final angular momenta of the atom.

In calculating the amplitudes of $E1$ transitions it is convenient to use the formalism of second quantization and the technique of summation over intermediate states.[10] This summation leads to the matrix element of a product of operators $D H_W \sim D(\mathbf{j}_e\mathbf{I})$ (here D is the operator for the electronic dipole moment and $H_W \sim \mathbf{j}_e\mathbf{I}$ according to Eq. (2), and can be calculated by expanding the product $D_i j_k$ into a sum of irreducible tensors:

$$D_i j_k = \sum_{s=0}^{2} T_{ik}^{s}$$

and using the formula[18]

$$\langle J'IF'\|[T^s \otimes I]\|JIF\rangle = \langle J'\|T^s\|J\rangle\langle I\|I\|I\rangle[3(2F+1)(2F'+1)]^{1/2}\begin{Bmatrix} J' & J & s \\ I & I & 1 \\ F' & F & 1 \end{Bmatrix},$$
$$\langle I\|I\|I\rangle = [I(I+1)(2I+1)]^{1/2}. \qquad (4)$$

The reduced matrix elements $\langle J'\|T^k\|J\rangle$ are calculated in the standard way.[10] The radial integrals needed for these calculations have been taken from[19] (see also[10,15]). We note that the technique for summing over intermediate states cannot be applied if the levels mixed by the weak interaction are close to each other. This sort of situation can occur in mixing of the state $6p^2 7s\,{}^4P_{1/2}'$ with $6p^3\,{}^2P_{1/2}$ and $6p^3\,{}^2P_{3/2}'$. In the former case the unusual closeness of the levels reduces the answer by 40 percent, but in the latter case the closeness is unimportant, since the matrix element

$$\langle 6p^2 7s\,{}^4P_{1/2}'|H_W|6p^3\,{}^2P_{1/2}'\rangle$$

is very small. We also note that because of strong can-

TABLE I. Amplitudes ($\langle J'IF'\|M1\|JIF\rangle/(-\mu_B)$) of $M1$ transitions and degrees of circular polarization in the transitions $6p_{1/2}$–$6p_{3/2}$, $\lambda = 12833$ Å in Tl and $6p^2\,{}^3P_0'$–$6p^2\,{}^3P_1$, $\lambda = 12789$ Å in Pb^{207}.

	F	F'	M1	$(P/\varkappa)\cdot 10^8$		F	F'	M1	$(P/\varkappa)\cdot 10^8$
Tl	0	1	−0.817	1.43	Pb	1/2	1/2	1.07	−0.150
	1	1	0.577	−0.475		1/2	3/2	−1.51	0.0748
	1	2	−1.29	−0.475					

TABLE II. Amplitudes ($\langle J'IF'\|M1\|JIF\rangle/(-\mu_B)$) of $M1$ transitions and degrees of circular polarization for transitions from the ground state to excited states in the configuration $6p^3$ of Bi.

F	F'	$^2D_{3/2}'$ (λ=8757 Å)		$^2D_{5/2}$ (λ=6477 Å)		$^2P_{1/2}'$ (λ=4616 Å)		$^2P_{3/2}'$ (λ=3015 Å)	
		M1	$(P/\varkappa)\cdot 10^8$	M1	$(P/\varkappa)\cdot 10^8$	M1	$(P/\varkappa)\cdot 10^8$	M1	$(P/\varkappa)\cdot 10^8$
3	2	—	—	−0.533	11.1	—	—	—	—
	3	1.53	−1.64	0.468	16.4	—	—	−0.169	3.13
	4	−1.69	−0.25	−0.306	23.5	−0.823	2.85	0.187	7.08
4	3	1.69	−4.09	−0.423	−0.797	—	—	−0.187	5.53
	4	0.15	19.5	0.568	6.29	−0.800	2.85	−0.0165	−84.9
	5	−1.95	2.17	−0.516	15.2	−0.482	−2.32	0.216	~1
5	4	1.95	−2.63	−0.309	−15.2	−0.652	2.85	−0.216	~−1
	5	−1.21	−2.62	0.570	−6.32	−0.800	−2.32	0.134	11.5
	6	−1.70	5.01	−0.719	4.31	—	—	0.189	−6.16
6	5	1.70	−0.75	−0.185	−32.1	−1.12	−2.32	−0.189	−8.49
	6	−2.60	1.34	0.471	−21.5	—	—	0.288	−2.56
	7	—	—	−0.923	−9.06	—	—	—	—

TABLE III. Maximum angles of rotation of the plane of polarization of the light, for Doppler shape of the absorption line of vapor at temperature 1200 °C.

	Transition	F	F'	$(\varphi/\varkappa)\cdot 10^8$, rad/m
Tl	$6p\ P_{1/2} - 6p\ P_{3/2}$	0	1	76
Pb	$6p^2\,{}^3P_0 - 6p^2\,{}^3P_1$	1/2	3/2	2.3
Bi	$6p^3\,{}^4S_{3/2}' - 6p^3\,{}^2D_{5/2}'$	5	6	2.7
	$^2D_{5/2}$	6	7	−1.4
	$^2P_{1/2}'$	6	5	−0.54
	$^2P_{3/2}'$	6	5	−0.06

TABLE IV. Amplitudes ($\langle J'IF' \| M1 \| JIF\rangle/(-\mu_B)$) and ($\langle J'IF' \| E1 \| JIF\rangle/(-\mu_B)$) of $M1$ and $E1$ transitions, and degree of polarization, for strongly forbidden transitions in Tl, Pb^{207}, and Cs.

	Transition	F	F'	$M1\cdot10^4$	$E1\cdot10^{12}$	$(P/\varkappa)\cdot10^5$
Tl	$6p_{1/2}-7p_{3/2}$ (λ=2927 Å)	0	1	0.366	7.48	−11.2
		1	0	−0.366	3.59	5.38
		1	1	0.517	−2.74	2.91
Pb	$6p^2\,{}^3P_0 \to 6p_{1/2}\,7p_{1/2}$ (λ=2330 Å)	1/2	1/2	–	1.65	—
		1/2	3/2	–	1.16	—
	$6p^2\,{}^3P_0 \to 6p_{1/2}\,7p_{3/2}$ (λ=2238 Å)	1/2	1/2	–	0.656	—
		1/2	3/2	–	0.464	—
Cs	$6s_{1/2}-7s_{1/2}$ (λ=5395 Å)	3	3	−0.973	−0.822	−0.463
		3	4	1.69	−1.43	0.465
		4	3	−1.69	−1.75	−0.568
		4	4	1.42	−0.932	0.360

cellation between the various contributions the accuracy of the calculation is very low for the transition $6p^3\,{}^4S_{3/2}' - 6p^3\,{}^2P_{3/2}'$ in bismuth.

Using the amplitudes for the $E1$ transitions found in this way, and amplitudes for the $M1$ transitions taken from[10,20], we can determine the degrees of circular polarization in the transitions in which we are interested. The values of $P/\varkappa$[see Eqs. (2) and (3)] and the amplitudes for the $M1$ transitions $6p_{1/2} - 6p_{3/2}$ in thallium, $6P^2\,{}^3P_0' - 6p^2\,{}^3P_1$ in lead, and $6p^3\,{}^4S_{3/2}' - 6p^3\,{}^2D_{3/2}'$, ${}^2D_{5/2}$, ${}^2P_{1/2}$, ${}^2P_{3/2}'$ in bismuth are shown in Tables I and II. The angles of rotation of the plane of polarization of the light from transitions between those hyperfine-structure components for which the rotation has maximum values are shown in Table III. All the values for degrees of circular polarization and angles of rotation are given without taking into account the contribution from the interaction between the electronic AC and the nucleonic VC, which was calculated previously in[10].

It can be seen from these results that to observe an effect of the type under consideration in bismuth the accuracy of measurement indicated in[5] would need to be raised by about one order of magnitude.

3. To complete the treatment of the problem, we calculate also the amplitudes for $E1$ transitions in the strongly forbidden $M1$ transitions $6s_{1/2} - 7s_{1/2}$ in cesium, $6p_{1/2} - 7p_{1/2}$ in thallium, and $6p^2\,{}^3P_0' - (6p_{1/2}7p_{1/2})$, $(6p_{1/2}7p_{3/2})$ in lead, which occur because of the interaction (1). Their values are shown in Table IV. For cesium and thallium, where the amplitudes of the $M1$ transitions are known from experiments,[21,3] this table also gives the degrees of circular polarization. The circulation polarization caused by the interaction of the electronic AC with the nucleonic VC in these transitions has been derived for cesium in[2,22] and for thallium and lead in[14,15]. We note that if there is no interaction between the electronic AC and the nucleonic VC, then it would probably be better to look for the circular polarization caused by the interaction between the electronic VC and the nucleonic AC in the transition $6p_{1/2} - 7p_{1/2}$ in thallium, rather than in 0–0 transitions in Pb^{207}.[14,15] The point is that although the value of $P/\varkappa$ is smaller for thallium by a factor three,[2)] the probability of the corresponding transition is almost three orders of magnitude larger than that of the 0–0 transition in Pb.[15]

The effects considered in this paper can be imitated by an external magnetic field. The restrictions on the magnetic field in the case of allowed $M1$ transitions can be obtained by using the results of[10,20]. For the most interesting transitions they are of the order of magnitude of 10^{-5} to 10^{-6} G.

[1)]The value of the constant $\varkappa g_I$ in thallium could probably be improved in accuracy if the contribution of the spin-orbit interaction to the configuration mixing in the ground state of the thallium nucleus were neglected. In this approximation the orbital angular momentum and the spin are separately conserved. Therefore the ground state can be characterized, despite the configuration mixing, by the values $L=0$, $S=\frac{1}{2}$ (in the shell model this means a proton in the ${}^2S_{1/2}$ state). Then from the experimentally known[17] magnetic moment of the thallium nucleus

$$(\mu_{Tl}=1,6=1/2 g_s{}^p\langle\sigma_z{}^p\rangle+1/2 g_s{}^n\langle\sigma_z{}^n\rangle,\quad \langle\sigma_z{}^n\rangle=1-\langle\sigma_z{}^p\rangle),$$

we find that $\langle\sigma_z p\rangle \approx \frac{3}{4}$, $\langle\sigma_z n\rangle \approx \frac{1}{4}$, i.e., $\varkappa_{Tl} \approx \frac{3}{4}\varkappa_p + \frac{1}{2}\varkappa_n$, and g_I is equal to two, as it is without taking any mixing into account.

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