Spectrum of electron-nuclear γ transitions in atomic nuclei

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The electronic state of an atom may change after emission or absorption of a γ quantum by the atomic nucleus, owing to the recoil of the nucleus and the inertia of the electronic shell. This leads to the appearance of electronic satellites in the nuclear γ-radiation spectrum. The positions and intensities of the electron-nuclear transition lines and means of their experimental observation are proposed. The shape of the Doppler contour is analyzed in the case when the satellites overlap the fundamental line because of the Doppler broadening. It is shown that when the atom is resonantly excited by coherent laser radiation the satellites are manifest as narrow resonances in the Doppler contour.

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1. INTRODUCTION. FORMULATION OF PROBLEM

The development of laser-spectroscopy methods has made it possible to register and use negligibly small changes, due to internal structure of the nucleus, in the structure of atomic and molecular spectra. Mention can be made here of effects such as the possibility of selective photoionization of atoms with an isomeric nucleus and the ensuing possibility of rapidly separating isomeric nuclei,[1] which is of interest for the γ-laser problem,[2-5] or the isomeric shift in the vibrational spectrum of a molecule, due to the increase of the excited-nucleus mass.[6]

On the other hand, the state of the atom or molecule should manifest itself also in quantum transitions, for example in the γ-ray spectrum of a nucleus. When a nucleus bound in an atom or molecule emits (absorbs) a γ ray, part of the nuclear energy can be transferred to (drawn from) the atom or molecule. In particular emission (absorption) of a γ ray by a nucleus in a molecule produces intense vibrational satellites in the γ-ray spectrum.[7] In the classical model, molecular-nuclear lines are due to modulation of the frequency of the nuclear emitter (absorber) by its vibrations in the molecule, and in the quantum model they are due to vibrational transitions of the molecule as it gives up momentum to the nucleus. The formation of molecular-nuclear lines is of interest because it uncovers a possibility of varying the γ-ray spectrum by changing the vibrational state of the molecule with coherent laser light.

A similar but smaller effect takes place also for nuclei in atoms. Within the framework of the classical model, the effect in atoms is due to the fact that the center of gravity of the nucleus does not coincide with the center of gravity of the entire atom, so that the momentum transferred to the nucleus influences the electron motion and vice versa. This gives rise to electronic satellites in the γ rays of the nucleus.[8] From very simple qualitative considerations we can also expect a certain acceleration of the "de-excitation" of the isomeric nuclei via transfer of part of the angular momentum to the electron shell, i.e., via radiative decay of metastable excited states of the nucleus, a decay accompanied by an electron-nuclear transition. The aim of this paper is to calculate the intensity and the structure of the electron-nuclear γ spectra in the atom, and also the effect of the electron shell on the decay rate of the metastable state of the nucleus.

It turns out that the intensity of the electron-nuclear satellites is fully sufficient to make them observable, and their spectrum can be controlled by exciting the atoms with laser radiation. However, owing to the smallness of the electron-proton interaction in comparison with the distances between the nuclear levels, the acceleration of the radiative de-excitation of the nucleus through the influence of the electron shell turns out to be quite insignificant. A possible exception are nuclear transitions between states with $J = 0$, for which one-photon transitions are forbidden.

2. GAMMA TRANSITIONS OF A NUCLEUS IN AN ATOM

Most excited states of nuclei have a many-particle character, and it is precisely with such states that one deals in experiment. Exceptions are the first excited states of nuclei with one or two nucleons or holes in excess of an even-even core.[9] The latter are more convenient for a theoretical analysis, since the single-particle model, to which we confine ourselves here, is applicable to them. The generalization of the formulas to the many-particle problem entails no difficulty and leads to no qualitatively new conclusions.

To investigate the influence of the electron shell on the nuclear transitions, we consider the aggregate of the protons and electrons that take part in the transition as a single quantum-mechanical system moving in the central field of a core whose state does not change in the course of the transition. We assume that only one proton and one electron take part in the transition. The Hamiltonian of the system is of the form

$$H=-\frac{\hbar^2}{2M}\left(\frac{\Delta_{\mathbf{r}_c}}{\mu_c}+\frac{\Delta_{\mathbf{r}_p}}{\mu_p}+\frac{\Delta_{\mathbf{r}_e}}{\mu_e}\right)+v_1(|\mathbf{r}_p-\mathbf{r}_c|)+v_2(|\mathbf{r}_e-\mathbf{r}_c|)-\frac{e^2}{|\mathbf{r}_p-\mathbf{r}_e|}; \tag{1}$$

here M is the mass of the entire atom, μ_c, μ_p, and μ_e are the relative masses of the core, proton, and electron ($\mu_x = m_x/M$), and $\Delta_{\mathbf{r}_c}$, $\Delta_{\mathbf{r}_p}$, and $\Delta_{\mathbf{r}_e}$ are the Laplace operators for the coordinates of the core, proton, and electron. As the zero-order approximation it is natural to consider the motion of a proton and electron that do not interact with each other in the field of the core. We make the following change of variables, which allows us to separate the variables in the zero-order Schrödinger equation:

$$\mathbf{R}=\mu_c\mathbf{r}_c+\mu_p\mathbf{r}_p+\mu_e\mathbf{r}_e, \tag{2}$$
$$\mathbf{R}_p=\mathbf{r}_p-\mathbf{r}_c, \quad \mathbf{R}_e=\mathbf{r}_e-\mathbf{r}_c.$$

The Hamiltonian of the system becomes in terms of the new variables

$$H=-\frac{\hbar^2}{2M}\left(\Delta+\frac{\mu_c+\mu_p}{\mu_c\mu_p}\Delta_p+\frac{\mu_e+\mu_c}{\mu_c\mu_e}\Delta_e+\frac{2\nabla_p\nabla_e}{\mu_c}\right)$$
$$+v_1(|\mathbf{R}_p|)+v_2(|\mathbf{R}_e|)-\frac{e^2}{|\mathbf{R}_p-\mathbf{R}_e|}, \qquad (3)$$

where Δ, Δ_p, and Δ_e are Laplace operators and ∇, ∇_p, ∇_e are the gradients with respect to the variables $\mathbf{R}$, $\mathbf{R}_p$, and $\mathbf{R}_e$, respectively. Separating the variables $\mathbf{R}$ from $\mathbf{R}_p$ and $\mathbf{R}_e$, we arrive at the problem of the motion of two particles (proton and electron) in a field:

$$v_1(|\mathbf{R}_p|)+v_2(|\mathbf{R}_e|)-\frac{e^2}{|\mathbf{R}_p-\mathbf{R}_e|}-\frac{\hbar^2\nabla_p\nabla_e}{2m_c}. \qquad (4)$$

The last term in (4) takes into account the dynamic interaction of two particles via a third one (the core).

In the zeroth approximation without allowance for the electron-proton interaction described by the last two terms in (4), the function of state of the system can be factorized:

$$\Psi(\mathbf{R},\mathbf{R}_p,\mathbf{R}_e)=\varphi_k(\mathbf{R})\psi_m(\mathbf{R}_p)\chi_s(\mathbf{R}_e), \qquad (5)$$

where $\varphi_k(\mathbf{R})$ is a function that describes the free motion of the entire atom, while $\psi_m(\mathbf{R}_p)$ and $\chi_s(\mathbf{R}_e)$ are functions of the bound states of the proton and electron in the central fields v_1 and v_2.

We consider now the system γ radiation accompanied by a change of the state of the nucleus and of the electron shell. The probability P of emitting a photon with momentum q and polarization e is proportional to the square of the modulus of the matrix element

$$\langle k'm's|V|kms\rangle, \qquad (6)$$

where V is the energy of the interaction of the photon with the atom:

$$V=N\left[\exp(-i\mathbf{qr}_e)\frac{e\nabla_{r_e}}{\mu_e}-\exp(-i\mathbf{qr}_p)\frac{e\nabla_{r_p}}{\mu_p}-Z\exp(-i\mathbf{qr}_c)\frac{e\nabla_{r_c}}{\mu_c}\right]; \qquad (7)$$

here Z is the charge of the core and N is the normalization of the photon function. In (7), the operator of the total current in the nucleus was replaced by the proton-current operator. According to a theorem by Siegert[10] this replacement is valid if $\mathbf{q}\cdot\mathbf{R}_{nuc}\ll 1$, which is usually satisfied. We change over in (7) to the coordinates $\mathbf{R}$, $\mathbf{R}_p$ and $\mathbf{R}_e$. Assuming that atom was at rest in the initial state, we leave out of V the term with ∇:

$$V=Ne^{-i\mathbf{qR}}\{\exp[-i\mathbf{q}(\mathbf{R}_e(\mu_c+\mu_e)-\mathbf{R}_p\mu_p)](e\nabla_e)\mu_c^{-1}$$
$$-\exp[-i\mathbf{q}(\mathbf{R}_p(\mu_c+\mu_p)-\mathbf{R}_e\mu_e)](e\nabla_p)\mu_p^{-1}+$$
$$+Z\exp[i\mathbf{q}(\mathbf{R}_p\mu_p+\mathbf{R}_e\mu_e)](\mathbf{e},\nabla_p+\nabla_e)\mu_c^{-1}\}. \qquad (8)$$

We neglect the contribution of the first term because of the rapid oscillations $\exp(-i\mathbf{q}\cdot\mathbf{R}_e)$. In the remaining terms we expand the exponentials in powers of q and confine ourselves only to the dipole proton and electron terms. This is justified by the smallness of $\mathbf{q}\cdot\mathbf{R}_p$ and $(\mathbf{q}\cdot\mathbf{R}_e)\mu_e$. In addition, we leave out terms that do not contain the proton coordinate, since we are not interested here in pure electronic transitions:

$$V_d\approx Ne^{-i\mathbf{qR}}\left[-\frac{e\nabla_p}{\mu_p}+i(\mathbf{qR}_e)(e\nabla_p)\frac{\mu_e}{\mu_p}+\right.$$
$$\left.+Z\frac{e\nabla_p}{\mu_c}+Zi(\mathbf{qR}_p)(e\nabla_e)\frac{\mu_p}{\mu_c}+Zi(\mathbf{qR}_e)(e\nabla_p)\frac{\mu_e}{\mu_c}\approx\right.$$
$$\left.\approx Ne^{-i\mathbf{qR}}\left[(e\nabla_p)\left(\frac{Z}{\mu_c}-\frac{1}{\mu_p}\right)+i(\mathbf{qR}_e)(e\nabla_p)\mu_e\left(\frac{Z}{\mu_c}+\frac{1}{\mu_p}\right)\right]. \qquad (9)$$

In the final expression for V_d we have neglected the term $Zi(\mathbf{q}\cdot\mathbf{R}_p)(\mathbf{e}\cdot\nabla_e)\mu_p/\mu_c$ in comparison with $Zi(\mathbf{q}\cdot\mathbf{R}_e)\times(\mathbf{e}\cdot\nabla_p)\mu_e/\mu_c$. The ratio of their matrix elements is of the order of $(R_a/R_{nuc})^2\mu_e/\mu_p$, where R_{nuc} and R_a are the radii of the nucleus and of the atom.

The first term in the brackets in expression (9) for the interaction energy leads to pure nuclear dipole transitions, while the second, which acts on the coordinates of the proton and the electron, is responsible for the electron satellites of the nuclear γ radiation. To take into account the nuclear transitions that are forbidden in the dipole approximation, it would be necessary to retain in (9) the next higher terms in the expansion in $\mathbf{q}\cdot\mathbf{R}_p$. We shall not do this, so as not to complicate the formulas, and note only that the analysis of the properties of the electron-nuclear transition spectra can be continued without specifying concretely the multipolarity of the nuclear transition—this has no effect whatever on the qualitative result.

The frequency of the γ transition of the nucleus with a change of the electronic state is given by

$$\hbar\omega_\gamma^{if}=\hbar\omega_{\gamma_0}\pm(\hbar\Delta_\gamma+E_i-E_f), \qquad (10)$$

where ω_{γ_0} is the frequency of the recoilless γ transition, $\hbar\Delta\gamma$ is the recoil energy, E_i and E_f are the initial and final energies of the electron; the plus sign corresponds to γ-quantum absorption and the minus sign to γ emission. If the atom was at the ground state ($E_i = 0$) prior to the emission of the γ ray, then all the electronic satellites are on the long-wave side (Fig. 1). Their distances from the ground-state line are equal to the excitation energies of the corresponding states of the electron shell. Analogous satellites, but on the short-wave side, should appear in the absorption spectra. The intensities of the various satellites, as follows from (9), are proportional to the squares of the dipole moments of the transitions between the initial and final states of the electron, and consequently decrease in proportion to n_f^{-3}, where n_f is the principal quantum number of the final state. The ratio of the intensities of the first satellite and the ground-state line can be estimated from the formula

$$k\approx(|\mathbf{q}_\gamma|\mathbf{R}_a\mu_e)^2. \qquad (11)$$

For example, for the 3-MeV transition in the isotope Ca^{41} [11] we have $k\approx 3\times10^{-4}$.

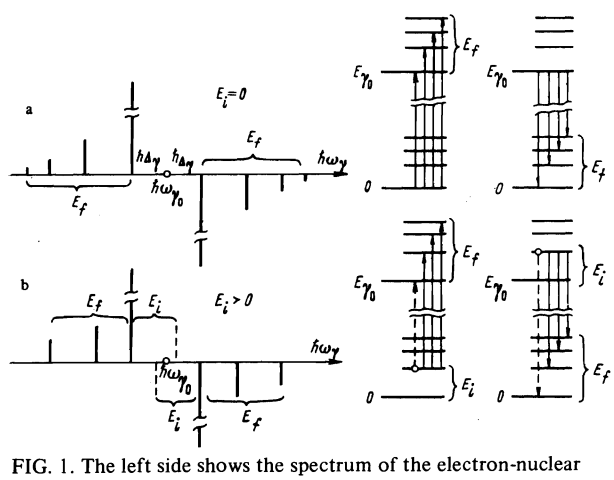

FIG. 1. The left side shows the spectrum of the electron-nuclear emission lines (the upward-directed lines) and absorption lines (downward) of the γ emission for a nucleus in the unexcited (a) and excited (b) atom. In the right-hand side are shown the corresponding quantum transitions in the "atom + nucleus" system (E_i and E_f are the initial and final energies of the electron, and E_{γ_0} is the energy of the nuclear level).

The appearance of electronic satellites is due to the inertia of the electron shell, which is excited by the rapid displacement of the nucleus as the result of the recoil. This interpretation of the phenomenon agrees with formula (11), which determines in fact the small parameter of the problem. For practical calculations it is convenient to use the following formula for the probability k_{0n} of exciting an atom to a level n following absorption of a γ photon (or of the inverse process):

$$k_{0n} \approx E_{0n}\,[\text{eV}]\ (f_{0n}E_{\gamma_0}\,[\text{MeV}]/A)^2, \tag{12}$$

where E_{0n} is the energy of the electronic transition in eV and f_{0n} is the oscillator strength of the 0–n transition.

Electronic satellites for unexcited atoms, removed from the central frequency $\omega_{\gamma 0}$, can be revealed by the overlap with the emission and absorption lines of other nuclei. Such an electronic rearrangement of the γ-transition frequency is significant if its magnitude exceeds the Doppler broadening of the γ line:

$$E_i - E_i \gg \hbar\omega_\gamma (2kT/Mc^2)^{1/2} = (4kTR)^{1/2}, \tag{13}$$

where $\hbar\Delta_\gamma = R = E_\gamma^2/2Mc^2$ is the recoil energy and T is the gas temperature. If relation (13) is not satisfied, then the excitation of the atom can be revealed by the ensuing fluorescence, which is perfectly observable. Thus, for an Ne^{21} nucleus absorbing a γ photon of energy $E_\gamma = 6$ MeV[11] formula (12) yields $k_{01} \approx 5 \times 10^{-3}$ for the probability of exciting the resonant Ne I level of energy $E_{01} = 16.7$ eV.

The excited atoms ($E_i > 0$) give electronic satellites that bring the absorption γ lines closer to the emission lines (dashed lines in Fig. 1). In this case the electronic satellite of the absorption (emission) γ line can be detected by coincidence with the ground-state emission (absorption) γ line. The condition for this is the equality $E_i - E_f = 2R$, which must be satisfied accurate to the γ-line Doppler broadening. This effect makes it possible to compensate in part for the detuning if the emission and absorption γ radiation in the gas phase. The methods used hitherto for this purpose were those of the fast relative motion of the source and target[12] and of heating the atoms.[13] The atoms can be excited in an electric discharge or with laser radiation. In the latter case it becomes possible in principle to obtain narrow γ emission (absorption) lines.

3. NARROW ELECTRON-NUCLEAR γ LINES INDUCED BY LASER RADIATION

The frequency of the γ transition of a nucleus, with allowance for the atom motion at a velocity $\mathbf{v}$, is given by Eq. (10) with a correction for the Doppler effect, namely $\hbar\omega_\gamma^{if} + \mathbf{q}_\gamma \cdot \mathbf{v}$. If the velocity distribution of the atoms is thermal (equilibrium), then the term $\mathbf{q}_\gamma \cdot \mathbf{v}$ describes the Doppler broadening of the γ line. However, when a plane coherent light wave acts coherently on low-pressure atomic vapor, one can observe atoms with definite velocity projections on the light-wave direction. We shall show that this makes it possible to obtain narrow frequency-tunable resonances of γ emission and absorption of nuclei in atoms, in the same manner proposed to produce this effect with nuclei in molecules.[7]

Let the coherent light wave act on a certain transition 0–n of the atom (Fig. 2). The action of the field causes the velocity distributions for the lower (0) and upper (n) levels to change in the following manner:

$$N_0(\mathbf{v}) = N_0^0(\mathbf{v})\frac{2+p(\mathbf{v})}{2[1+p(\mathbf{v})]},$$

$$N_n(\mathbf{v}) = N_n^0(\mathbf{v})\frac{p(\mathbf{v})}{2[1+p(\mathbf{v})]}, \tag{14}$$

where $N_0^0(\mathbf{v})$ and $N_n^0(\mathbf{v})$ are the equilibrium distributions of the particle velocities at the levels 0 and n, and the parameter $p(\mathbf{v})$ describes the dependence of the degree of saturation of the absorption of the moving atom on the detuning of its transition frequency ω_{0n} relative to the light-field frequency ω:

$$p(v) = \frac{G\Gamma^2}{(\omega - \omega_{0n} - \mathbf{kv})^2 + \Gamma^2}. \tag{15}$$

Under the condition of exact resonance ($\omega - \omega_{0n} = \mathbf{k}\cdot\mathbf{v}$), the parameter $p(\mathbf{v})$ coincides with the degree of saturation G of the transition. The relative widths of the spectral intervals inside the Doppler-broadened optical and nuclear lines, which correspond to atoms (and the nuclei in them) interacting with the coherent wave, are equal to each other:

$$\Gamma_{\text{nuc}}/\omega_{\gamma_0} = \Gamma(1+G)^{1/2}/\omega_{0n}. \tag{16}$$

The shape of the spectral absorption (emission) line of the electronic satellite of the γ transition at the frequency ω_γ^{if}, due to the nuclei in the atoms that are situated at the i-th level (i = 0, n) at the initial instant of time and at the level f after the absorption (emission) of the γ quantum, is described by the expression

$$F_{if}(\omega_\gamma) = \int \sigma(\omega_\gamma - \omega_\gamma^{if} - \mathbf{q}_\gamma\mathbf{v}) N_i(\mathbf{v})\, d(\mathbf{v}\mathbf{n}_\gamma), \tag{17}$$

where $\mathbf{n}_\gamma$ is a unit vector in the direction of the γ quantum, the distributions $N_i(\mathbf{v})$ are given by (14), $\sigma(x)$ is the cross section for resonant absorption of a γ quantum by a free nucleus and is described by the expression

$$\sigma(x) = \sigma_0\Gamma_\gamma^2/(x^2 + \Gamma_\gamma^2), \tag{18}$$

where σ_0 is the absorption cross section at the maximum and Γ_γ is the natural half-width of the γ transition of the nucleus.

If the γ radiation is observed in the same direction as the light wave ($\mathbf{n}_\gamma = \mathbf{n}$), then a narrow structure is produced in the spectrum $F_{if}(\omega\gamma)$ and is due to a hole in the velocity distribution $N_0(v_z)$ or to a peak in the distribution $N_n(v_z)$, where $v_z = \mathbf{v}\cdot\mathbf{n}$ is the projection of the molecule velocity on the direction of the light wave ($\mathbf{n}$ and $\mathbf{n}_\gamma$ are directed along the z axis). If the relative natural width of the γ transition $\Gamma_\gamma/\omega_{\gamma_0}$ is much smaller than the relative homogeneous width of the electronic transition, i.e.,

$$\Gamma_\gamma/\omega_{\gamma_0} \ll (\Gamma/\omega_0)(1+G)^{1/2}, \tag{19}$$

then the width $2\Gamma_{\text{nuc}}$ of the produced narrow γ resonances is due to an optical "channel" and is determined by expression (16). In this case the integral (17) reduces to the simple form:

$$\Gamma_{if}(\omega_\gamma) = \sigma_0 N_i(\xi), \quad \xi = (\omega_\gamma - \omega_\gamma^{if})/q_\gamma, \tag{20}$$

where $N_i(\xi) = N_i(v_z)$ are the functions (14) that describe the distributions of the projections of the atom velocities on the direction of the light wave at each of the initial levels of the atom.

We consider first excited-atom nuclei that have electronic satellites at the frequencies ω_γ^{nf}. The summary form of the spectral line of the γ transition for these nuclei is given by

$$\Gamma_{nf}(\omega_\gamma) = \sigma_0 \sum_j k_{nj} N_n\left(\frac{\omega_\gamma - \omega_\gamma^{nf}}{q_\gamma}\right), \tag{21}$$

where $\omega_\gamma^{nf} = \omega_{\gamma 0} \pm (\Delta_\gamma + \omega_{nf})$ and $\hbar\omega_{nf} = E_f - E_n$. It follows from (21) and (14) that at the resonant frequencies

$$\omega_{res}^{nf} = \omega_\tau^{nf} + (\omega_{\tau_0}/\omega_{0n})\Omega, \quad \Omega = \omega - \omega_{0n}, \tag{22}$$

there are produced narrow peaks with half-width Γ_n determined by (16) and with relative amplitude described by (11). A typical form of the spectrum $F_{nf}(\omega\gamma)$ is shown in Fig. 3a.

The nuclei in the unexcited atoms have electronic satellites at the frequencies ω_γ^{0f}, and the summary form of the spectral line of the γ transition for these nuclei is given by

$$F_{0f}(\omega_\tau) = \sigma_0 \sum_f k_{0f} N_0 \left(\frac{\omega_\tau - \omega_\tau^{0f}}{q_\tau}\right), \tag{23}$$

where $\omega_\gamma^{0f} = \omega_{\gamma 0} \pm (\Delta_\gamma + \omega_{0f})$ and $\hbar\omega_{0f} = E_f$. In this case there are produced at the resonant frequencies

$$\omega_{res}^{0f} = \omega_\tau^{0f} + (\omega_{\tau_0}/\omega_{0n})\Omega \tag{24}$$

narrow dips with half-width Γ_{nuc} and relative amplitude described by the relation (11). A typical form of the spectrum $F_{nf}(\omega\gamma)$ is shown in Fig. 3b.

The observed quantity is obviously the summary spectrum

$$F(\omega_\gamma) = F_{0f}(\omega_\gamma) + F_{nf}(\omega_\gamma), \tag{25}$$

due to the nuclei in the unexcited and excited atoms. We consider first the central component formed by the lines $F_{00}(\omega\gamma)$ and $F_{nn}(\omega\gamma)$. The frequencies of the dip in the $F_{00}(\omega\gamma)$ line and of the peak in $F_{nn}(\omega\gamma)$ coincide, and the amplitudes of the central components are the same ($k_{00} = k_{nn} \approx 1$). Since the depth of the hole in the distribution $N_0(v_z)$ is exactly equal to the amplitude of the peak in $N_n(v_z)$, we obtain

$$F_{00}(\omega_\tau) + F_{nn}(\omega_\tau) = \sigma_0 N\left(\frac{\omega_{\tau_0} \pm \Delta_\tau}{q_\tau}\right), \tag{26}$$

where $N(\xi) = N_0^0(\xi) + N_n^0(\xi)$, i.e., it represents an ordinary structureless Doppler contour.

At the frequency ω_γ^{n0} of the "red" electronic satellite of the absorption line there is produced only a narrow peak, the frequency of which ω_{res}^{n0} can be tuned within the limits of the Doppler contour of the γ-transition satellite when the light-wave frequency is scanned along the Doppler contour of the 0–n optical transition. Similar resonant absorption peaks of width $2\Gamma_{nuc}$ are produced at the frequencies $\omega_\gamma^{n,f>n}$ of the "blue" electronic satellites that are separated from the central component by the frequency of the f–n electronic transition. They can, however, become superimposed on the Doppler contours of the "blue" satellites at the frequencies ω_γ^{0f}, which in turn contain narrow resonant dips of width $2\Gamma_{nuc}$ at the frequencies ω_{res}^{0f}, which are also tunable as the light-wave frequency is scanned.

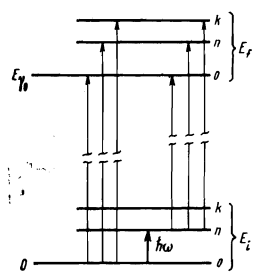

FIG. 2. Scheme of electron-nuclear quantum transitions that participate in the formation of the γ absorption lines of the nucleus in an atom in which laser radiation acts on the 0–n transition.

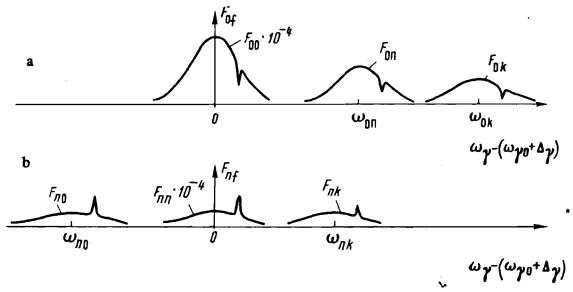

FIG. 3. Narrow resonances inside Doppler-broadened electron-nuclear γ-absorption lines for nuclei in atoms at the lower level (a) and in excited atoms (b).

4. INFLUENCE OF ELECTRON SHELL ON THE TOTAL PROBABILITY

We consider in first-order perturbation theory the effect of the electron-proton interaction described by the last two terms of (4) on the total probability of the γ decay of the nucleus, i.e., for example, the possibility of accelerating the radiative de-excitation of a metastable nucleus through interaction with the electron shell.

The interaction $\nabla_p \nabla_e / m_c$ mixes the proton states with orbital momentum values L that differ by unity, and consequently it is possible to lower by unity the multipolarity of the allowed transition. Let us estimate the small parameter α of this interaction. The usual formulas of perturbation theory lead to the expression

$$\alpha = \hbar^2 \frac{(\langle\nabla_p\rangle, \langle\nabla_e\rangle)}{m_c E_p} \approx 2\left(\frac{m_p m_e}{m_c^2}\right)^{1/2}\left(\frac{E_e}{E_p}\right)^{1/2} \approx \frac{1}{20A}\left(\frac{E_e}{E_p}\right)^{1/2}, \tag{27}$$

where E_e and E_p are the binding energies of the electron in the atom and of the proton in the nucleus. For the transition considered above in Ne^{21} we have $\alpha^2 \approx 10^{-9}$, whereas in an isolated nucleus the transition is attenuated by a factor 10^{-4} to 10^{-5} when the multipolarity is increased by unity.

The higher harmonics of the Coulomb electron-proton interaction lead to a mixing of the porton states with $\Delta L \leq j$, where j is the number of the harmonic. The small parameter is equal in this case to

$$\alpha \approx \frac{e^2}{R_a E_p}\left(\frac{R_{nuc}}{R_a}\right)^j. \tag{28}$$

For $R_a = 1$Å and $E_p = 3$ MeV the parameter is $\alpha^2 \approx 10^{-12}$ $(5 \times 10^{-5})^{2j}$. We see that in this case, too, the probability of the 2^{k-j}-pole nuclear transitions induced is much smaller than that of the 2^k-pole transitions that are allowed in an isolated nucleus.

Thus, our analysis shows that in those cases when some single-quantum transitions are allowed in the nucleus, their probability always exceeds the probability of the transitions induced by the electron shell.

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