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INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

ON QUANTUM QUADRUPOLE RADIATION

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ON QUANTUM QUADRUPOLE RADIATION*

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ABSTRACT

In this paper it is shown that for the electromagnetic decay of a quantum system in a coherent rotational state the total quadrupole radiation is proportional to $\overset{\cdot\cdot\cdot\cdot\cdot}{Q}\,\ddot{Q}^{*}$ + c.c. For the radiation flux out of a sphere of large radius a different quantity, closer to the classical expression $\dddot{Q}^2$, is found.

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1. Introduction

We study in this paper the quadrupole radiation power for the electromagnetic decay of coherent rotational states. The γ- decay of these states has been considered in previous papers with particular attention to the pulsed behaviour of their time evolution.(1) Within the framework of ref. 1, we evaluate in Sect. 2.a the total quadrupole radiation power and find that it is proportional to $\overset{\cdot\cdot\cdot\cdot\cdot}{Q}\,\ddot{Q}^{*}$ + c.c.. In Sect. 2.b the power emitted out of a sphere of very large radius is evaluated and found to be proportional to $|\dddot{Q}^{(+)}|^2$, where $Q^{(+)}$ is the operator obtained from Q via suppression of the off-diagonal elements $\langle Q\rangle_{l'l}$ with $l'<l$. This result is therefore very close to the classical expression $\dddot{Q}^{2}$ (2).

We shall consider the electromagnetic decay of a coherent rotational state (*)

$$|\psi_\alpha(t)\rangle = \sum_{I} a_I^\alpha\, e^{-iE_I t}|\phi_I\rangle \tag{1.1}$$

which has at time $t = 0$ the following characteristic properties: a) the absolute values of the amplitudes a_I^α are peaked around a mean value of the angular momentum I, b) the phases of a_I^α are roughly equidistant, c) the energies E_I of the states ϕ_I obey the rule $\omega_0 I(I+1)$ to a very good approximation. The states ϕ_I are eigenstates of H_0 defined as the sum of the nuclear and the free electromagnetic Hamiltonians. The total Hamiltonian of the system is then given by

$$\begin{aligned} H &= H_0 + H_I \\ H_0 &= H_{nuclear} + H^0_{e.m.} \end{aligned} \tag{1.2}$$

(*) Here and in what follows the index $\mathcal{I}$ stands for both angular momentum quantum numbers I and M. We use natural units $\hbar = c = 1$.

where H_I is the quadrupole interaction:

$$(1.3) \qquad H_I = \frac{1}{6(2\pi)^{3/2}} \int d^3k \sqrt{\frac{k}{2}} \left\{ a_k^\dagger \left[Q_k, H_0 \right] + h.c. \right\}$$

By the symbol of integration we mean both the integration over momenta $\vec{k}$ and the sum over polarization vectors $\vec{e}_k$. Throughout the paper, whenever k appears as a subscript it denotes both momentum $\vec{k}$ and polarization $\vec{e}_k$ of the photon. A cut-off for large k will be always understood in the integrals over the variable k, in order to be consistent with $kR_m \ll 1$, where R_m is the radius of the nucleus.

The quadrupole moment Q_k is given by

$$(1.4) \qquad Q_k = \sum_{i=1}^{Z} 3e \frac{(\vec{k}\cdot\vec{r}_i)(\vec{e}_k\cdot\vec{r}_i)}{k}$$

In order to get (1.3) we have assumed that the nuclear force is velocity independent and we have dropped the magnetic part of the interaction, which is supposed to be small.

2. Evaluation of radiation power

We shall denote by Ψ_α the initial coherent rotational state of our decaying system. One and two-photon states will be denoted by $\Psi_{\beta k}$ and $\Psi_{\beta k k'}$, respectively. Since we shall use perturbation theory in the interaction picture, we shall write the evolution of these states via the unperturbed Hamiltonian H_0:

$$(2.1) \qquad \begin{aligned} |\Psi_\alpha(t)\rangle &\equiv e^{-iH_0 t}|\Psi_\alpha\rangle = \sum_I a_I^\alpha e^{-iE_I t}|\phi_I\rangle \\ |\Psi_{\beta k}(t)\rangle &\equiv a_k^\dagger e^{-ikt}|\Psi_\beta(t)\rangle \\ |\Psi_{\beta k k'}(t)\rangle &\equiv a_k^\dagger a_{k'}^\dagger e^{-i(k+k')t}|\Psi_\beta(t)\rangle \end{aligned}$$

We shall not consider the contribution of states having more than two photons. The solution of the complete Schrödinger equation:

$$(2.2) \qquad H|\Psi(t)\rangle = i\frac{d|\Psi\rangle}{dt}$$

is then expanded on a complete orthonormal set of 0, 1 and 2 photon states:

$$(2.3) \qquad \begin{aligned} |\Psi(t)\rangle = &\sum_\beta c_\beta(t)|\Psi_\beta(t)\rangle + \sum_\beta \int d^3k\, c_{\beta k}(t)|\Psi_{\beta k}(t)\rangle + \\ &+ \sum_\beta \iint d^3k\, d^3k'\, c_{\beta k k'}(t)|\Psi_{\beta k k'}(t)\rangle \end{aligned}$$

For the coefficients, to second order, we get:

$$(2.4) \qquad c_\beta(t) = \delta_{\beta\alpha} - i\sum_{\beta'} \int_0^t dt' \int d^3k\, c_{\beta' k}(t')\, e^{-ikt'} \langle \Psi_\beta(t')| H_I a_k^\dagger |\Psi_{\beta'}(t')\rangle$$

$$(2.5) \qquad c_{\beta k}(t) = -i\sum_{\beta'} \int_0^t dt'\, c_{\beta'}(t')\, e^{ikt'} \langle \Psi_\beta(t')| a_k H_I |\Psi_{\beta'}(t')\rangle$$

$$(2.6) \qquad c_{\beta k k'}(t) = -i\sum_{\beta'} \int_0^t dt' \int d^3k''\, c_{\beta' k''}(t')\, e^{i(k+k'-k'')t'} \langle \Psi_\beta(t')| a_k a_{k'} H_I a_{k''}^\dagger |\Psi_{\beta'}(t')\rangle$$

2.a. Emission out of a sphere of nuclear radius

Conservation of energy tells us that

$$
\begin{aligned}
(2.7)\qquad & \frac{d}{dt}\langle\Psi(t)|\int\mathcal{H}\,dv\,|\Psi(t)\rangle \equiv \\
& \equiv \frac{d}{dt}\langle\Psi(t)|\left\{\int\mathcal{H}_{nuclear}\,dv+\int\mathcal{H}^{\circ}_{e.m.}\,dv+\int\mathcal{H}_{I}\,dv\right\}|\Psi(t)\rangle = 0
\end{aligned}
$$

where $\mathcal{H}$, $\mathcal{H}_{nuclear}$, $\mathcal{H}^{\circ}_{e.m.}$ and $\mathcal{H}_I$ are the Hamiltonian densities corresponding to H, $H_{nuclear}$, $H_{e.m.}$ and H_I respectively. In (2.7) the integrations are of course performed over all coordinate space. Since $\mathcal{H}_{nuclear}$ and $\mathcal{H}_I$ are non-vanishing only within the small sphere of nuclear radius R_n, from (2.7) we get:

$$
\begin{aligned}
(2.8)\qquad & \frac{d}{dt}\langle\Psi(t)|\int_{\text{within } R_n}\mathcal{H}\,dv\,|\Psi(t)\rangle = \\
& = -\frac{d}{dt}\langle\Psi(t)|\int_{\text{outside } R_n}\mathcal{H}^{\circ}_{e.m.}\,dv\,|\Psi(t)\rangle \cong -\frac{d}{dt}\langle\Psi(t)|\int_{\text{all space}}\mathcal{H}^{\circ}_{e.m.}\,dv\,|\Psi(t)\rangle
\end{aligned}
$$

that is, the change in energy of the system within the sphere of nuclear radius R_n manifests itself as the change of energy of the free electromagnetic field in all space.

Evaluation of the right hand side of (2.8) yields:

$$
\begin{aligned}
(2.9)\qquad P(t) & \equiv \frac{d}{dt}\langle\Psi(t)|\int_{\text{all space}}\mathcal{H}^{\circ}_{e.m.}\,dv\,|\Psi(t)\rangle = \\
& = \frac{d}{dt}\langle\Psi(t)|\int d^3k\; k\, a^{+}_{k}a_{k}\,|\Psi(t)\rangle
\end{aligned}
$$

which, to second order, turns out to be

$$
(2.10)\qquad P(t) = \frac{d}{dt}\sum_{\beta}\int d^3k\;|c_{\beta k}(t)|^2\,k
$$

In order to evaluate this quantity, we follow just the same procedure used in Ref. 1 for the evaluation of $W_{\alpha\to\beta}$. One finally gets:

$$
\begin{aligned}
(2.11)\qquad P(t) & = \frac{d}{dt}\frac{1}{36\pi}\sum_{\beta}\sum_{ll'LL'}\left\{\Delta^5_{ll'}\Delta_{LL'}\Theta(\Delta_{ll'})+\Delta^5_{LL'}\Delta_{ll'}\Theta(\Delta_{LL'})\right\}\cdot \\
& \cdot\overline{Q_{l'l}Q^{*}_{L'L}}\;\frac{\sin(\Delta_{ll'}-\Delta_{LL'})t/2}{\Delta_{ll'}-\Delta_{LL'}}\;e^{i(\Delta_{LL'}-\Delta_{ll'})t/2} = \\
& = \frac{1}{72\pi}\sum_{\beta}\sum_{ll'LL'}\left\{\Delta^5_{ll'}\Delta_{LL'}\Theta(\Delta_{ll'})+\Delta^5_{LL'}\Delta_{ll'}\Theta(\Delta_{LL'})\right\}\cdot \\
& \cdot\overline{Q_{l'l}Q^{*}_{L'L}}\;e^{i(\Delta_{LL'}-\Delta_{ll'})t}
\end{aligned}
$$

where by $\overline{QQ^{*}}$ we mean that the average over directions and polarizations has been performed. We have defined $Q_{l'l}$ and $\Delta_{l'l}$ as

$$
(2.12)\qquad
\begin{aligned}
Q_{l'l} & = a^{\beta *}_{l'}a^{\alpha}_{l}\langle\phi_{l'}|Q|\phi_{l}\rangle \\
\Delta_{l'l} & = E_{l'}-E_{l}
\end{aligned}
$$

Since the initial state is supposed to be a coherent rotational state, the contribution coming from the initial state α is the dominant one. We drop then the sum on β and retain only the $\beta=\alpha$ term. One gets a simpler form for (2.11) by exchanging $l'\leftrightarrow L$ and $l\leftrightarrow L'$ in the sum pertaining to the second term in the curly bracket and using the fact that for $\beta=\alpha$ one has $Q_{l'l}=Q^{*}_{ll'}$. Using the fact that $\Theta(x) + \Theta(-x) = 1$, one finally gets

$$(2.13)\quad P(t) = \frac{1}{72\pi} \sum_{ll'LL'} \Delta^5_{ll'} \Delta_{LL'} \overline{Q_{l'l} Q^*_{L'L}}\, e^{i(\Delta_{LL'} - \Delta_{ll'})t}$$

Expression (2.13) can be obtained by taking the average over directions and polarizations of the following quantity:

$$(2.14)\quad P(t) = \frac{1}{12^2\pi} \overline{\ddddot{Q}_\alpha(t)\, \dot{Q}^*_\alpha(t)} + c.c.$$

where $Q_\alpha(t)$ is defined as the mean value of the quadrupole operator for the coherent rotational state $\Psi_\alpha(t)$,

$$(2.15)\quad Q_\alpha(t) = \langle \Psi_\alpha(t) | Q | \Psi_\alpha(t) \rangle$$

2.b. Emission out of a sphere of large radius

We shall now evaluate the radiation power emitted from a sphere of large (but finite) radius R . Eq. (2.8) will read now

$$(2.16)\quad \frac{d}{dt} \langle \Psi(t) | \int_{\text{within } R} \mathcal{H}\, dv | \Psi(t) \rangle = -\frac{d}{dt} \langle \Psi(t) | \int_{\text{outside } R} \mathcal{H}^0_{e.m.}\, dv | \Psi(t) \rangle$$

Using at the right-hand side of (2.16) the continuity equation for the photon field, one gets for the radiation power:

$$(2.17)\quad P_R(t) = -\langle \Psi(t) | \int_{S_R} \vec{J} \cdot \vec{n}\, ds | \Psi(t) \rangle$$

where S_R is the surface of the sphere of radius R and $\vec{J}$ is the Poynting vector:

$$(2.18)\quad \vec{J} = -N \left[\frac{\partial \vec{A}}{\partial t} \times \mathrm{rot}\, \vec{A} \right]$$

Here N means normal product for the creation and annihilation photon operators.

For $\vec{J}$ we easily get

$$(2.19)\quad \vec{J}(\vec{r}) = \vec{J}_0(\vec{r}) + \vec{J}_2(\vec{r})$$

where

$$(2.20)\quad \vec{J}_0(\vec{r}) = (2\pi)^{-3} \iint \frac{d^3k\, d^3k'}{2} \sqrt{kk'}\; \vec{e}_k \times \left(\frac{\vec{k}'}{k'} \times \vec{e}_{k'} \right) \cdot \left(a^+_{k'} a_k\, e^{i(\vec{k} - \vec{k}')\cdot\vec{r}} + h.c. \right)$$

$$(2.21)\quad \vec{J}_2(\vec{r}) = (2\pi)^{-3} \iint \frac{d^3k\, d^3k'}{2} \sqrt{kk'}\; \vec{e}_k \times \left(\frac{\vec{k}'}{k'} \times \vec{e}_{k'} \right) \cdot \left(-a_k a_{k'}\, e^{i(\vec{k} + \vec{k}')\cdot\vec{r}} + h.c. \right)$$

The operator $\vec{J}_0$ does not change the number of photons in the field, while $\vec{J}_2$ changes that number by two.

We shall first show that $\vec{J}_2$ does not give contribution to the emitted power. Let us take its average on the state $|\Psi(t)\rangle$ and use (2.3) and (2.4-6):

$$\vec{J}_2(\vec{r},t) \equiv \langle \Psi(t)|\, \vec{J}_2(\vec{r})\, |\Psi(t)\rangle = \tag{2.22}$$

$$= (2\pi)^{-3} \sum_{\beta} \iint d^3k\, d^3k'\, \sqrt{kk'}\; \vec{e}_k \times \left(\frac{\vec{k}'}{k'} \times \vec{e}_{k'} \right) \cdot$$

$$\cdot\, c_\beta^*(t)\, c_{\beta k k'}(t)\, e^{i(\vec{k}+\vec{k}')\cdot\vec{r}}\, e^{-i(k+k')t}$$

Since $c_{\beta k k'}$ is at least $O(H_I^2)$, c_β must be taken to the zeroth order, i.e. $c_\beta(t) \cong \delta_{\beta\alpha}$. Then, apart from a constant, for $\vec{J}_2(\vec{r},t)$ one gets

$$\vec{J}_2(\vec{r},t) = \sum_{\beta} \sum_{11'LL'} \iint d^3k\, d^3k'\; kk'\, \vec{e}_k \times \left(\frac{\vec{k}'}{k'} \times \vec{e}_{k'} \right) \cdot$$

$$\cdot\, e^{i(\vec{k}+\vec{k}')\vec{r}}\, e^{-i(k+k')t}\, a_{1'}^{\alpha *}\, a_1^{\beta}\, a_{L'}^{\beta *}\, a_L^{\alpha}\; \Delta_{LL'}\, \Delta_{11'} \cdot \tag{2.23}$$

$$\cdot \left\{ \langle Q_k \rangle_{L'L} \langle Q_{k'} \rangle_{1'1}\, \frac{1}{k-\Delta_{LL'}} \cdot \left[\frac{e^{i(k+k'-\Delta_{11'}-\Delta_{LL'})t} - 1}{k+k'-\Delta_{11'}-\Delta_{LL'}} - \right.\right.$$

$$\left.\left. - \frac{e^{i(k'-\Delta_{11'})t} - 1}{k'-\Delta_{11'}} \right] + (\vec{k} \leftrightarrow \vec{k}') \right\} + c.c.$$

We evaluate now (2.23) as follows: If $\Delta_{LL'}$ is negative, we do not have any contribution to the integral (this is of course approximate). The same holds for $\Delta_{LL'} = 0$, since $\Delta_{LL'}$ appears as a factor in the integral. If $\Delta_{LL'}$ is positive, we add the interval $(-\infty, 0)$ and evaluate the integral by contour integration in the k complex plane (this is possible since a cut-off function is understood in the integral). With these approximations, one just gets the contributions at the poles $k = \Delta_{LL'}$ and $k = \Delta_{LL'} + \Delta_{11'} - k'$, which are both zero:

$$\vec{J}_2(\vec{r},t) \cong 0 \tag{2.24}$$

Let us now come to the evaluation of $\vec{J}_0$. Its average on the state $|\Psi(t)\rangle$ is given by

$$\vec{J}_0(\vec{r},t) \equiv \langle \Psi(t)|\, \vec{J}_0(\vec{r})\, |\Psi(t)\rangle =$$

$$= (2\pi)^{-3} \sum_{\beta} \iint d^3k\, d^3k'\, \sqrt{kk'}\; \vec{e}_k \times \left(\frac{\vec{k}'}{k'} \times \vec{e}_{k'} \right) \cdot$$

$$\cdot\, \mathrm{Re}\left[c_{\beta k'}^*(t)\, c_{\beta k}(t)\, e^{i(\vec{k}-\vec{k}')\cdot\vec{r}}\, e^{i(k'-k)t} \right]$$

Using (2.5) and (2.17) one finally gets:

$$P_R(t) = -\frac{R^2}{36\,(2\pi)^6} \sum_{\beta} \sum_{11'LL'} a_{1'}^{\alpha *}\, a_1^{\beta}\, a_{L'}^{\beta *}\, a_L^{\alpha}\, e^{-\frac{i}{2}(\Delta_{11'}+\Delta_{LL'})t} \cdot$$

$$\cdot\, \Delta_{11'}\, \Delta_{LL'} \mathop{\iint\!\!\!\!\!\!\!\sum}_{pp'} d^3k\, d^3k'\; kk'\, \langle Q_{\vec{k}'p'} \rangle_{1'1} \langle Q_{\vec{k}p} \rangle_{L'L}\, e^{\frac{i}{2}(k'-k)t} \cdot \tag{2.25}$$

$$\cdot\, \frac{\sin(k'+\Delta_{11'})t/2}{k'+\Delta_{11'}}\, \frac{\sin(k-\Delta_{LL'})t/2}{k-\Delta_{LL'}}\, \mathcal{J}_R(\vec{k}p, \vec{k}'p') + c.c.$$

where $\mathcal{J}_R(\vec{\kappa}p, \vec{\kappa}'p')$ is given by:

$$\mathcal{J}_R(\vec{\kappa}p, \vec{\kappa}'p') = \int d\Omega_R \left[\vec{e}_{\kappa p} \times \left(\frac{\vec{\kappa}'}{\kappa'} \times \vec{e}_{\kappa' p'}\right)\right] \cdot \vec{n}\, e^{i(\vec{\kappa}-\vec{\kappa}')\cdot\vec{R}} \tag{2.26}$$

where we have explicitely shown the sums over polarizations p and p' .
In order to proceed with the calculation one takes advantage of the slow dependence of the quadrupole matrix elements on polarizations and angular directions. One then extracts these matrix elements from the integrals and sums by taking their average values $\overline{\langle Q \rangle_{I'I}} \cdot \overline{\langle Q \rangle_{L'L}}$.
The evaluation of the spatial angular integral, summed over polarizations and integrated over angles in $\vec{K}$ and $\vec{K}'$ momentum spaces, for R large (remember the presence of cut-offs in the K and K' integrations) gives:

$$\int d\Omega_\kappa \int d\Omega_{\kappa'} \sum_{pp'} \int d\Omega_R\, \mathcal{J}_R(\vec{\kappa}p, \vec{\kappa}'p') \underset{R\ large}{\sim} \underset{R\ large}{\sim} \frac{(4\pi)^3}{2\kappa\kappa' R^2}\, e^{i\kappa R} e^{-i\kappa' R} \tag{2.27}$$

where, by the argument of the stationary phase (note that $t > 0$), we have dropped all terms of the type $\exp ih(\pm R \pm \frac{t}{2})$, where h here stands for either k or k' , which give vanishing small contributions when integrated over k or k' . In order to integrate on k and k', one

can follow the same analytic procedure used for the evaluation of $\vec{J}_2$.
We get:

$$P_R(t) = -\frac{1}{72\pi} \sum_\beta \sum_{II'LL'} a_{I'}^{\alpha *} a_I^\beta a_{L'}^{\beta *} a_L^\alpha\, \Delta_{II'}^3 \Delta_{LL'}^3 \cdot \Theta(-\Delta_{II'})\, \Theta(\Delta_{LL'})\, \overline{\langle Q \rangle_{I'I}}\, \overline{\langle Q \rangle_{L'L}}\, e^{i(\Delta_{II'}+\Delta_{LL'})(R-t)} + c.c. \tag{2.28}$$

Again, since the initial state is supposed to be a coherent rotational state, the contribution of the initial state α is dominant; we drop then the sum on β and retain only the $\beta=\alpha$ term. It is then immediately seen that the complex conjugate term equals the first term at the right hand side of (2.28) (to see this interchange $I \leftrightarrow L'$ and $L \leftrightarrow I'$ in the c.c. term). If we define an operator $Q^{(+)}$ such that

$$\langle Q^{(+)} \rangle_{I'I} = \langle Q \rangle_{I'I}\, \Theta(\Delta_{I'I}) \tag{2.29}$$

we can finally write the radiation power out of a large sphere of radius R as:

$$P_R(t) = \frac{1}{36\pi} \left| \dddot{Q}_\alpha^{(+)}(t-R) \right|^2 \tag{2.30}$$

where $Q_\alpha^{(+)}(t-R)$ is given by

$$Q_\alpha^{(+)}(t-R) = \langle \Psi_\alpha(t-R) | Q^{(+)} | \Psi_\alpha(t-R) \rangle \tag{2.31}$$

and it is obtained from Q_α (as defined by (2.15)) by suppression of the off-diagonal terms $\langle Q \rangle_{I'I}$ with $I' < I$. The result (2.30) is very close to the classical expression $\frac{1}{180} \dddot{Q}^2$ for the radiation power(2).

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