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MODIFYING THE LIFETIME OF AN UNSTABLE SYSTEM BY AN INTENSE ELECTROMAGNETIC FIELD*

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We study the temporal behavior of a three-level system (such as an atom or a molecule), initially prepared in an excited state, bathed in a laser field tuned at the transition frequency of the other level. We analyze the dependence of the lifetime of the initial state on the intensity of the laser field. The phenomenon we discuss is related to both electromagnetic induced transparency and quantum Zeno effect.

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1 Introduction

The temporal behavior of quantum mechanical systems, being governed by unitary operators [1], displays some subtle features at short [2] and long times [3]. In order to discuss the evolution of genuine unstable systems one usually makes use of the Weisskopf-Wigner approximation [4], which ascribes the main properties of the decay law to a pole located near the real axis of the complex energy plane. This yields the Fermi "Golden Rule" [5]. These features of the quantum evolution are well known and discussed in textbooks of quantum mechanics [6] and quantum field theory [7]. For a recent review, see [8]. In this paper we shall investigate the possibility that the lifetime of an unstable quantum system can be modified by the presence of a very intense electromagnetic field. We shall look at the temporal behavior of a three-level system (such as an atom or a molecule), where level #1 is the ground state and levels #2, #3 are two excited states. The system is initially prepared in level #2 and if it follows its natural evolution, it will decay to level #1. The decay will be (approximately) exponential and characterized by a certain lifetime, that can be calculated from the Fermi Golden Rule. But what happens if one shines on the system an intense laser field, tuned at the transition frequency 3-1? This problem was investigated in Ref. [9], in the context of the so-called quantum Zeno

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effect [10]. It was found that the lifetime of the initial state depends on the intensity of the laser field. In the limit of extremely intense field, the decay should be considerably slowed down. The aim of this paper is to study this effect in more detail and discuss a new phenomenon: we shall see that for physically sensible values of the laser field, the decay can be *enhanced*, rather than hindered. This can be viewed as an "inverse" quantum Zeno effect. The whole problem is related to electromagnetic induced transparency [11].

2 Preliminaries and definitions

We start from the Hamiltonian [9]:

$$\begin{aligned} H &= H_0 + H_{\rm int} \\ &= \omega_0|2\rangle\langle 2| + \Omega_0|3\rangle\langle 3| + \sum_{\boldsymbol{k},\lambda}\omega_k a^\dagger_{\boldsymbol{k}\lambda}a_{\boldsymbol{k}\lambda} + \sum_{\boldsymbol{k},\lambda}\left(\phi_{\boldsymbol{k}\lambda}a^\dagger_{\boldsymbol{k}\lambda}|1\rangle\langle 2| + \phi^*_{\boldsymbol{k}\lambda}a_{\boldsymbol{k}\lambda}|2\rangle\langle 1|\right) \\ &\quad + \sum_{\boldsymbol{k},\lambda}\left(\Phi_{\boldsymbol{k}\lambda}a^\dagger_{\boldsymbol{k}\lambda}|1\rangle\langle 3| + \Phi^*_{\boldsymbol{k}\lambda}a_{\boldsymbol{k}\lambda}|3\rangle\langle 1|\right), \end{aligned} \tag{1}$$

where the first two terms are the free Hamiltonian of the 3-level system (whose states $|i\rangle$ $(i = 1, 2, 3)$ have energies $E_1 = 0$, $\omega_0 = E_2 - E_1 > 0$, $\Omega_0 = E_3 - E_1 > 0$), the third term is the free Hamiltonian of the EM field and the last two terms describe the $1 \leftrightarrow 2$ and $1 \leftrightarrow 3$ transitions in the rotating wave approximation, respectively. The states $|2\rangle$ and $|3\rangle$ are chosen so that no transition between them is possible (e.g., because of selection rules). The matrix elements of the interaction Hamiltonian read

$$\begin{aligned} \phi_{\boldsymbol{k}\lambda} &= \frac{e}{\sqrt{2\epsilon_0 V\omega}}\int d^3x\; e^{-i\boldsymbol{k}\cdot\boldsymbol{x}}\boldsymbol{\epsilon}^*_{\boldsymbol{k}\lambda}\cdot\boldsymbol{j}_{12}(\boldsymbol{x}), \\ \Phi_{\boldsymbol{k}\lambda} &= \frac{e}{\sqrt{2\epsilon_0 V\omega}}\int d^3x\; e^{-i\boldsymbol{k}\cdot\boldsymbol{x}}\boldsymbol{\epsilon}^*_{\boldsymbol{k}\lambda}\cdot\boldsymbol{j}_{13}(\boldsymbol{x}), \end{aligned} \tag{2}$$

where $-e$ is the electron charge, ϵ_0 the vacuum permittivity, V the volume of the box, $\omega = |\boldsymbol{k}|$, $\boldsymbol{\epsilon}_{\boldsymbol{k}\lambda}$ the photon polarization and $\boldsymbol{j}_{fi}$ the transition current of the radiant system. For example, in the case of an electron moving in an external field, we have $\boldsymbol{j}_{fi} = \psi^\dagger_f\boldsymbol{\alpha}\psi_i$ where ψ_i and ψ_f are the electron wavefunctions of the initial and final states, respectively, and the components of $\boldsymbol{\alpha}$ are the usual Dirac matrices. For the sake of generality we are using relativistic matrix elements (although our analysis can be performed with nonrelativistic ones $\boldsymbol{j}_{fi} = \psi^*_f\boldsymbol{p}\psi_i/m_e$, where $\boldsymbol{p}/m_e$ is the electron velocity).

We shall concentrate our attention on a 3-level system bathed in a continuous laser beam, whose photons have momentum $\boldsymbol{k}_0$ ($|\boldsymbol{k}_0| = \Omega_0$) and polarization λ_0. We shall also assume, throughout this paper, that

$$\phi_{\boldsymbol{k}_0\lambda_0} = 0, \tag{3}$$

i.e., the laser does not interact with state $|2\rangle$. Also, since the average number N_0 of $\boldsymbol{k}_0$-photons in the total volume V can be considered very large, we shall perform our analysis in terms of number (rather than coherent) states of the EM field [12]. In this approximation,

$$\begin{aligned}\langle 1; 0_{\boldsymbol{k}\lambda}, N_0 | H_{\rm int} | 3; 0_{\boldsymbol{k}\lambda}, N_0 - 1\rangle &= \sqrt{N_0}\,\Phi_{\boldsymbol{k}_0\lambda_0} \\ &\gg \langle 1; 1_{\boldsymbol{k}\lambda}, N_0 - 1 | H_{\rm int} | 3; 0_{\boldsymbol{k}\lambda}, N_0 - 1\rangle = \Phi_{\boldsymbol{k}\lambda},\end{aligned} \tag{4}$$

$\forall(\boldsymbol{k},\lambda) \neq (\boldsymbol{k}_0,\lambda_0)$. In the above equation and henceforth, the vector $|i; n_{\boldsymbol{k}\lambda}, M_0\rangle$ represents an atom or a molecule in state $|i\rangle$, with $n_{\boldsymbol{k}\lambda}$ $(\boldsymbol{k},\lambda)$-photons and M_0 laser photons.

In the above approximation, the Hamiltonian (1) can be replaced by

$$\begin{aligned} H \;\simeq\; & \omega_0|2\rangle\langle 2| + \Omega_0|3\rangle\langle 3| + \sum_{\boldsymbol{k},\lambda}\omega_k a^\dagger_{\boldsymbol{k}\lambda} a_{\boldsymbol{k}\lambda} + \sum_{\boldsymbol{k},\lambda}{}' \left(\phi_{\boldsymbol{k}\lambda} a^\dagger_{\boldsymbol{k}\lambda}|1\rangle\langle 2| + \phi^*_{\boldsymbol{k}\lambda} a_{\boldsymbol{k}\lambda}|2\rangle\langle 1|\right) \\ & + \left(\Phi_{\boldsymbol{k}_0\lambda_0} a^\dagger_{\boldsymbol{k}_0\lambda_0}|1\rangle\langle 3| + \Phi^*_{\boldsymbol{k}_0\lambda_0} a_{\boldsymbol{k}_0\lambda_0}|3\rangle\langle 1|\right), \end{aligned} \tag{5}$$

where the prime means that the summation does not include $(\boldsymbol{k}_0,\lambda_0)$ [due to our hypothesis (3)]. The operators

$$\mathcal{N} = |2\rangle\langle 2| + \sum_{\boldsymbol{k},\lambda}{}' a^\dagger_{\boldsymbol{k}\lambda} a_{\boldsymbol{k}\lambda}, \qquad \mathcal{N}_0 = |3\rangle\langle 3| + a^\dagger_{\boldsymbol{k}_0\lambda_0} a_{\boldsymbol{k}_0\lambda_0} \tag{6}$$

satisfy

$$[H,\mathcal{N}] = [H,\mathcal{N}_0] = [\mathcal{N}_0,\mathcal{N}] = 0, \tag{7}$$

which imply the conservation of the total number of photons plus the atomic excitation (Tamm-Dancoff approximation [13]). The Hilbert space splits therefore into sectors that are invariant under the action of the Hamiltonian: in our case, the system evolves in the subspace labelled by the eigenvalues $\mathcal{N} = 1, \mathcal{N}_0 = N_0$ and the analysis can be restricted to this sector [12].

3 Temporal evolution

The states of the total system in the sector $(\mathcal{N},\mathcal{N}_0) = (1, N_0)$ read

$$|\psi(t)\rangle = x(t)|2; 0, N_0\rangle + \sum_{\boldsymbol{k},\lambda}{}' y_{\boldsymbol{k}\lambda}(t)|1; 1_{\boldsymbol{k}\lambda}, N_0\rangle + \sum_{\boldsymbol{k},\lambda}{}' z_{\boldsymbol{k}\lambda}(t)|3; 1_{\boldsymbol{k}\lambda}, N_0 - 1\rangle, \tag{8}$$

with the normalization

$$\langle\psi(t)|\psi(t)\rangle = |x(t)|^2 + \sum_{\boldsymbol{k},\lambda}{}' |y_{\boldsymbol{k},\lambda}(t)|^2 + \sum_{\boldsymbol{k},\lambda}{}' |z_{\boldsymbol{k},\lambda}(t)|^2 = 1, \qquad \forall t. \tag{9}$$

At time $t = 0$ we prepare our system in the state

$$|\psi(0)\rangle = |2; 0, N_0\rangle \quad \Leftrightarrow \quad x(0) = 1,\ y_{\boldsymbol{k}\lambda}(0) = 0,\ z_{\boldsymbol{k}\lambda}(0) = 0, \tag{10}$$

which is an eigenstate of the free Hamiltonian

$$H_0|\psi(0)\rangle = H_0|2; 0, N_0\rangle = \omega_0|2; 0, N_0\rangle. \tag{11}$$

We set, without any loss of generality, $E_1 + N_0\Omega_0 = 0$. By inserting (8) in the time-dependent Schrödinger equation

$$i\frac{d}{dt}|\psi(t)\rangle = H|\psi(t)\rangle \tag{12}$$

and Laplace transforming with the initial conditions (10), one readily obtains

$$\widetilde{x}(s) = \frac{1}{s + i\omega_0 + Q(B,s)}, \tag{13}$$

$$\widetilde{y}_{\boldsymbol{k}\lambda}(s) = \frac{-i\phi_{\boldsymbol{k}\lambda}(s+i\omega_k)}{(s+i\omega_k)^2 + B^2}\,\widetilde{x}(s), \tag{14}$$

$$\widetilde{z}_{\boldsymbol{k}\lambda}(s) = -\frac{\sqrt{N_0}\Phi^*_{\boldsymbol{k}_0\lambda_0}\phi_{\boldsymbol{k}\lambda}}{(s+i\omega_k)^2 + B^2}\,\widetilde{x}(s), \tag{15}$$

where the tilde denotes Laplace transform and

$$Q(B,s) = \sum_{\boldsymbol{k},\lambda} |\phi_{\boldsymbol{k}\lambda}|^2 \frac{s+i\omega_k}{(s+i\omega_k)^2 + B^2}, \qquad B^2 = N_0\,|\Phi_{\boldsymbol{k}_0\lambda_0}|^2. \tag{16}$$

B is the intensity of the laser field and can be viewed as the “strength” of the observation performed by the laser beam on level #2 [9].

In the continuum limit ($V \to \infty$), the matrix elements scale as follows

$$\lim_{V\to\infty} \frac{V\omega^2}{(2\pi)^3} \sum_{\lambda} \int d\Omega |\phi_{\boldsymbol{k}\lambda}|^2 \equiv g^2\omega_0\chi^2(\omega), \tag{17}$$

where Ω is the solid angle. The (dimensionless) function $\chi(\omega)$ and the coupling constant g have the following general properties:

$$\chi^2(\omega) \propto \begin{cases} \omega^{2j\mp1} & \text{if } \omega \ll \Lambda \\ \omega^{-\beta} & \text{if } \omega \gg \Lambda \end{cases}, \tag{18}$$

$$g^2 = \alpha(\omega_0/\Lambda)^{2j+1\mp1}, \tag{19}$$

where j is the total angular momentum of the photon emitted in the $2 \to 1$ transition, $\mp$ represent electric and magnetic transitions, respectively, $\beta(>1)$ is a constant, α the fine structure constant and Λ a natural cutoff (e.g., of the order of the inverse Bohr radius), which determines the range of the atomic or molecular form factor. The above equations are due to general properties of the matrix elements [14, 15].

In order to scale the quantity B, we take the limit of very large cavity, by keeping the density of Ω_0-photons in the cavity constant:

$$V \to \infty, \qquad N_0 \to \infty, \quad \text{with} \quad \frac{N_0}{V} = n_0 = \text{const.} \tag{20}$$

We obtain from (16)

$$B^2 = n_0 V |\Phi_{\boldsymbol{k}_0\lambda_0}|^2 = (2\pi)^3 n_0 |\varphi_{\lambda_0}(\boldsymbol{k}_0)|^2 \tag{21}$$

where φ is the scaled matrix element of the 1-3 transition. As we shall see, in order to affect significantly the lifetime of level #2, we shall need a high value of B. It is therefore of interest to consider a 1-3 transition of the dipole type, in which case the above formula reduces to

$$B^2 = 2\pi\alpha\Omega_0|\boldsymbol{\epsilon}^*_{\boldsymbol{k}_0\lambda_0} \cdot \boldsymbol{x}_{13}|^2 n_0, \tag{22}$$

where $\boldsymbol{x}_{13}$ is the dipole matrix element.

3.1 Laser off

Let us first look at the case $B = 0$. The laser, tuned at the 1-3 frequency Ω_0, is off and we expect to recover the well-known physics of a two-level system prepared in an excited state and coupled to the radiation field [16]. In this case, $Q(0, s)$ is nothing but the self-energy function

$$Q(s) \equiv Q(0, s) = \sum_{\boldsymbol{k},\lambda} |\phi_{\boldsymbol{k}\lambda}|^2 \frac{1}{s + i\omega_k} \tag{23}$$

and becomes, in the continuum limit,

$$Q(s) \equiv g^2\omega_0 q(s) \equiv -ig^2\omega_0 \int_0^\infty d\omega \frac{\chi^2(\omega)}{\omega - is}, \tag{24}$$

where χ is defined in (17). The function $\widetilde{x}(s)$ in Eq. (13) (with $B = 0$) has a logarithmic branch cut extending from 0 to $-i\infty$, no singularities on the first Riemann sheet (physical sheet) and a simple pole on the second Riemann sheet. The pole equation is

$$s + i\omega_0 + g^2\omega_0 q_{\rm II}(s) = 0, \tag{25}$$

where

$$q_{\rm II}(s) = q(se^{-2\pi i}) = q(s) + 2\pi\chi^2(is) \tag{26}$$

is the determination of $q(s)$ on the second Riemann sheet. We note that $g^2q(s)$ is $O(g^2)$, so that the pole can be found perturbatively: by expanding $q_{\rm II}(s)$ around $-i\omega_0$ we get a power series, whose radius of convergence is $R_c = \omega_0$ because of the branch point at the origin. The circle of convergence lies half on the first Riemann sheet and half on the second sheet (Fig. 1). The pole is well inside the convergence circle, because $|s_{\rm pole} + i\omega_0| \sim g^2\omega_0 \ll R_c$, and we can write

$$s_{\rm pole} = -i\omega_0 - g^2\omega_0 q_{\rm II}(-i\omega_0 - 0^+) + O(g^4) = -i\omega_0 - g^2\omega_0 q(-i\omega_0 + 0^+) + O(g^4), \tag{27}$$

because $q_{\rm II}(s)$ is the analytical continuation of $q(s)$ below the branch cut. By using the formula

$$\lim_{\varepsilon\to 0^+} \frac{1}{x \pm i\varepsilon} = P\frac{1}{x} \mp i\pi\delta(x), \tag{28}$$

one gets from (24)

$$\begin{aligned} q(-i\eta + 0^+) &= -i\int_0^\infty d\omega\, \chi^2(\omega) \frac{1}{\omega - \eta - i0^+} \\ &= \pi\chi^2(\eta)\theta(\eta) - iP\int_0^\infty d\omega\, \chi^2(\omega)\frac{1}{\omega - \eta} \end{aligned} \tag{29}$$

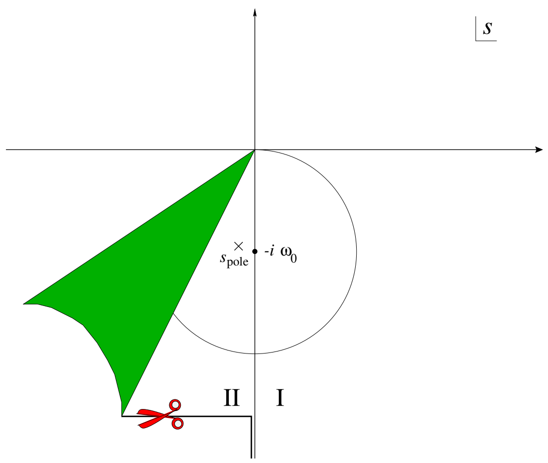

Fig. 1. Cut and pole in the s-plane ($B = 0$). I and II are the first and second Riemann sheets, respectively.

and by setting

$$s_{\rm pole} = -i\omega_0 + i\Delta E - \frac{\gamma}{2}, \tag{30}$$

one gets

$$\gamma = 2\pi g^2\omega_0\chi^2(\omega_0) + O(g^4), \qquad \Delta E = g^2\omega_0 P\int_0^\infty \frac{\chi^2(\omega)}{\omega-\omega_0} + O(g^4), \tag{31}$$

which are the Fermi Golden Rule and the second order correction to the energy ω_0 of level #2.

3.2 Laser on

We now turn our attention to the situation with the laser switched on, $B \neq 0$. The self energy function $Q(B,s)$ in (16) depends on the value of B and can be written in terms of the self energy function $Q(s)$ in absence of laser field [Eq. (23)], by making use of the following remarkable property:

$$\begin{aligned} Q(B,s) &= \frac{1}{2}\sum_{\boldsymbol{k},\lambda}|\phi_{\boldsymbol{k}\lambda}|^2\left(\frac{1}{s+i\omega_k+iB}+\frac{1}{s+i\omega_k-iB}\right) \\ &= \frac{1}{2}\left[Q(s+iB)+Q(s-iB)\right]. \end{aligned} \tag{32}$$

Notice, incidentally, that in the continuum limit ($V \to \infty$), due to the above formula, $Q(B,s)$ scales like $Q(s)$. The position of the pole $s_{\rm pole}$ (and as a consequence the lifetime $\tau_{\rm E} \equiv \gamma^{-1} = -1/2\Re e s_{\rm pole}$) depends on the value of B. There are now two branch cuts in the complex s plane, due to the two terms in (32). They lie over the imaginary axis, along $(-i\infty, -iB]$ and $(-i\infty, +iB]$.

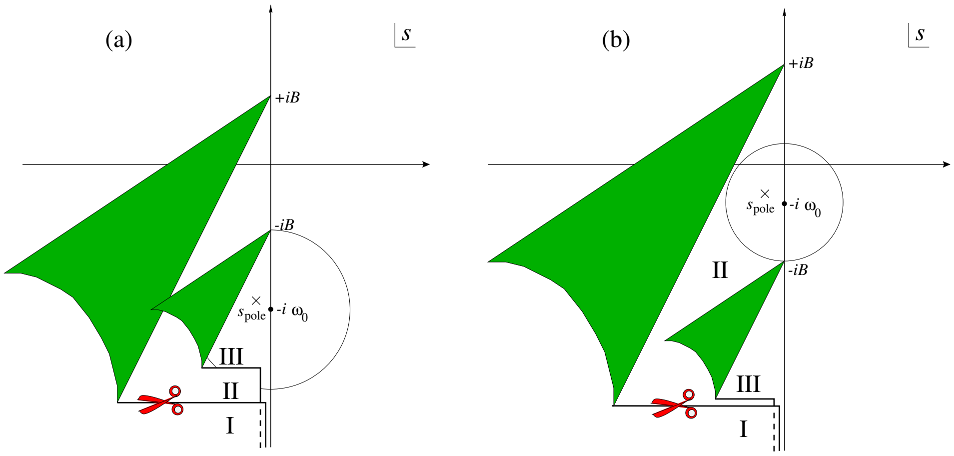

Fig. 2. Cuts and pole in the s-plane ($B \neq 0$) and convergence circle for the expansion of $Q(B, s)$ around $s = -i\omega_0$. I , II and III are the first, second and third Riemann sheets, respectively. (a) $B < \omega_0$. (b) $B > \omega_0$.

The pole satisfies the equation

$$s + i\omega_0 + Q(B, s) = 0, \tag{33}$$

where $Q(B, s)$ is of order g^2, as before, and can again be expanded in power series around $s = -i\omega_0$, in order to find the pole perturbatively. However, this time one has to choose the right determination for the function $Q(B, s)$. There are two cases: a) The branch point $-iB$ is situated above $-i\omega_0$, so that $-i\omega_0$ lies on both cuts. See Fig. 2(a); b) The branch point $-iB$ is situated below $-i\omega_0$, so that $-i\omega_0$ lies only on the upper branch cut. See Fig. 2(b). In case a), i.e. for $B < \omega_0$, the pole is on the third Riemann sheet (under both cuts) and the power series converges in a circle lying half on the first and half on the third Riemann sheet, within a convergence radius $R_c = \omega_0 - B$, which decreases as B increases [Fig. 2(a)]. In case b), i.e. for $B > \omega_0$, the pole is on the second Riemann sheet (under the upper cut only) and the power series converges in a circle lying half on the first and half on the second Riemann sheet, within a convergence radius $R_c = B - \omega_0$, which increases with B [Fig. 2(b)]. In both cases we can write, for $|s_{\rm pole} + i\omega_0| < R_c = |B - \omega_0|$,

$$\begin{aligned} s_{\rm pole} &= -i\omega_0 - \frac{1}{2}\left\{Q[-i(\omega_0 + B) + 0^+] + Q[-i(\omega_0 - B) + 0^+]\right\} + O(g^4) \\ &= -i\omega_0 - \frac{1}{2}g^2\omega_0\left\{q[-i(\omega_0 + B) + 0^+] + q[-i(\omega_0 - B) + 0^+]\right\} + O(g^4). \end{aligned} \tag{34}$$

Our analysis is therefore valid when $|B - \omega_0|$ is larger than $g^2\omega_0$, (we remind the reader that we are interested in large values of B). Equation (34) enables us to analyze the behavior of the lifetime of level #2.

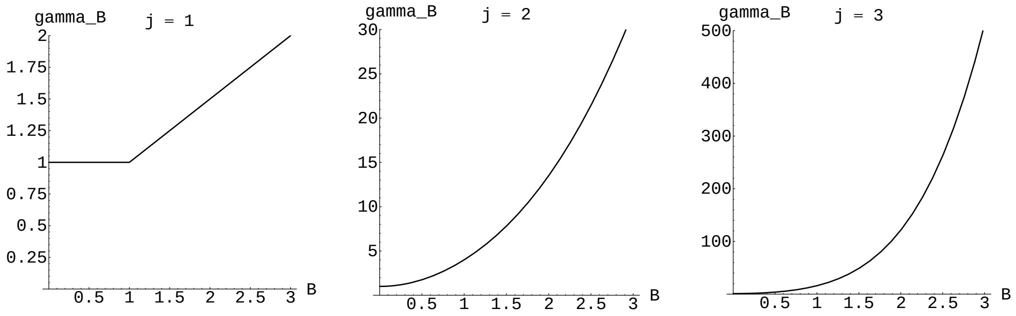

Fig. 3. The decay rate $\gamma(B)$ vs B, for electric transitions with $j = 1, 2, 3$; $\gamma(B)$ is in units γ and B in units ω_0. Notice the different scales on the vertical axis.

4 Decay rate vs B

We write, as in (30), $s_{\text{pole}} = -i\omega_0 + i\Delta E(B) - \gamma(B)/2$. By substituting (29) into (34) and taking the real part, one obtains the following expression for the decay rate

$$\gamma(B) = \pi g^2 \omega_0 \left[\chi^2(\omega_0 + B) + \chi^2(\omega_0 - B)\theta(\omega_0 - B)\right] + O(g^4), \tag{35}$$

which becomes, taking into account (31),

$$\gamma(B) = \frac{\gamma}{2} \frac{\chi^2(\omega_0 + B) + \chi^2(\omega_0 - B)\theta(\omega_0 - B)}{\chi^2(\omega_0)} + O(g^4). \tag{36}$$

This is our main result and involves no approximations. It expresses the lifetime $\gamma(B)^{-1}$, when the system is bathed in an intense laser field B, in terms of the "ordinary" lifetime γ^{-1}, when there is no laser field. By taking into account the general behavior (18) of the matrix elements $\chi^2(\omega)$ and substituting into (36), one gets to $O(g^4)$

$$\gamma(B) \simeq \frac{\gamma}{2} \left[\left(1 + \frac{B}{\omega_0}\right)^{2j\mp 1} + \left(1 - \frac{B}{\omega_0}\right)^{2j\mp 1} \theta(\omega_0 - B)\right], \qquad (B \ll \Lambda) \tag{37}$$

where $\mp$ refers to 1-2 transitions of electric and magnetic type, respectively. Observe that, since $\Lambda \sim$ inverse Bohr radius, the case $B \ll \Lambda$ is the physically relevant one [15]. The decay rate is profoundly modified by the presence of the "B"-field. Its behavior is shown in Fig. 3 for a few values of j. The case $j = 1$ (transition of electric dipole type) yields a constant value up to $B = \omega_0$; above this threshold, γ increases linearly with B. For $j > 1$ the derivative of $\gamma(B)$ is continuous. In general, the decay rate $\gamma(B)$ increases with B, so that the lifetime $\gamma(B)^{-1}$ decreases as B is increased. If one looks at B as the strength of the "observation" performed by the laser beam on level #2 [9], one can view this phenomenon as an "inverse" quantum Zeno effect.

Equation (37) is valid for $B \ll \Lambda$. In the opposite case, $B \gg \Lambda$, by (18) and (36), one gets to $O(g^4)$

$$\gamma(B) \sim \frac{\gamma}{2} \frac{\chi^2(B)}{\chi^2(\omega_0)} \propto B^{-\beta}. \qquad (B \gg \Lambda) \tag{38}$$

This is essentially the result obtained in Ref. [9]. If such high values of B were experimentally obtainable, the decay would be considerably hindered (quantum Zeno effect [9]). However, in such a case, ionization effects should be considered and our model should be refined in order to take them into account. Similar physical limitations were also emphasized by Kurizki, in a slightly different context [17].

A thorough investigation of the phenomenon we have just proposed is in preparation [15]. It involves a complete discussion in terms of Fano states and electromagnetic induced transparency [11] and an analysis of the decay rates from level #2 to the dressed atomic or molecular states.

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