Quantum averaging and resonances: two-level atom in a one-mode classical laser field

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Abstract
We use a nonperturbative method based on quantum averaging and an adapted from of resonant transformations to treat the resonances of the Hamiltonian of a two-level atom interacting with a one-mode classical field in Floquet formalism. We illustrate this method by extraction of effective Hamiltonians of the system in two regimes of weak and strong coupling. The results obtained in the strong-coupling regime, are valid in the whole range of the coupling constant for the one-photon zero-field resonance.

Keywords: quantum averaging, two-level atom, Floquet formalism, resonant transformation, nonlinear resonance

1. Introduction
Treatment of resonances in quantum systems cannot be performed by perturbative approaches. In the vicinity of resonances the perturbative formulas display small denominators that lead to the divergence of the perturbative expansions. A widely use model that incorporates a one-photon resonance is the Jaynes-Cummings Hamiltonian extracted from the full dressed Hamiltonian that describes a two-level system coupled with a single mode of a quantized field [1]. Its counterpart for an interaction with a classical laser field in the rotating-wave approximation (RWA) is the RWA Hamiltonian [2]. The RWA is valid only when an atom is perturbed by a weak and near-resonant (classical or quantized) field. For large values of the field amplitude (or coupling constant), the RWA cannot be invoked and one has to take in to account additionally the counter-rotating terms of the Hamiltonian. The goal is to obtain the spectrum for a whole interval of values of a parameter like the coupling constant beyond RWA. This is needed e.g. in applications where the coupling changes adiabatically [3], corresponding e.g. to envelopes of laser pulses or to transversal spatial profiles of cavity fields.

The semiclassical model of a two-level atom coupled with several incommensurate frequencies [4] of a classical field has been treated by different methods in Refs. [5-8]. In Ref. [9], we used the method of quantum averaging and resonant transformations proposed in [10] to detect the resonances of the model of a two-level atom interacting with a single-mode quantized field in two regimes of weak and strong coupling.

In this paper, we use the same method as [10] with an adapted form of resonant transformation to obtain the effective Hamiltonians of the model of a two-level atom interacting with a single-mod classical (laser) field in Floquet formalism. In the weak-coupling regime we have to iterate this procedure several times to derive the essential structure of the spectrum in large ranges of the coupling constant. In the strong-coupling regimes, the qualitative properties of the spectrum can be globally obtained by some preliminary unitary transformations and one resonant transformation which treat the zero-field resonances. In the case of zero-field one-photon resonance we obtain an accurate approximation, valid for all values of the coupling constant, that contains all the qualitative structure. Once this main structure is obtained, one can systematically improve the quantitative accuracy of the spectrum by applying perturbative methods.

This paper is organized as follows: In the next section, we explain the quantum averaging technique, the resonance concept, KAM-type perturbative transformations, and resonant transformation. In Sec. 3 we overview the Floquet formalism, to be self-contained, which is an essential tool to study laser-driven systems. Section 4 is devoted to the presentation of the model and some preliminary considerations. In Sec. 5, taking into account the resonances of this model in the weak-coupling regime, we extract the effective Hamiltonians by quantum averaging technique and resonant transformations. In Sec. 6 we extract the effective Hamiltonians in the strong-coupling regime. Finally, in Sec. 7 we give some conclusions and summary.
2. Resonance analysis and the method

We consider a Hamiltonian \( H = H_0 + \epsilon V \) where \( H_0 \) is the reference (unperturbed) Hamiltonian, \( \epsilon V \) is the perturbation and \( \epsilon \) is an ordering parameter [18]. The first analysis of this problem is in terms of perturbation theory: we look for a KAM-type unitary transformation \( e^{iW_1} \) close to the identity that allows reducing the order of the perturbation from \( \epsilon \) to \( \epsilon^2 \):

\[
e^{-\epsilon W_1} H e^{\epsilon W_1} = H_0 + \epsilon D_1 + \epsilon^2 V_2. \tag{1}
\]

\( \epsilon D_1 \) is a remaining term of order \( \epsilon \) that satisfies \( [H_0, D_1] = 0 \). The unknown \( W_1 \) and \( D_1 \) are solution of the following equations (Cohomological equation) [10, 11]:

\[
[H_0, W_1] + V_1 = D_1, \tag{2a}
\]

\[
[H_0, D_1] = 0. \tag{2b}
\]

The remaining perturbation of order \( \epsilon^2 \) is given by

\[
\epsilon^2 V_2 = \sum_{m=2}^{\infty} \frac{\epsilon^m}{m!} (m-1)_m^{-1} W_1 + \frac{\epsilon^m}{m!} D_1, \tag{3}
\]

where \( W_1 \) is defined as:

\[
W_1 = [B, W_1]. \tag{4}
\]

The solutions of eqs. (2a,b) can be written in terms of averaging [10,12-14]:

\[
D_1 = \mathcal{P} = \prod_{H_0} V := \lim_{\tau \to +\infty} \frac{1}{\tau} \int_0^\tau ds e^{-iH_0 S} V e^{iH_0 S} = \sum_{\nu, j, j'} \langle \nu, j | V | \nu, j' \rangle \langle \nu, j' |,
\tag{5a}
\]

\[
W_1 = \lim_{\tau \to +\infty} \frac{1}{\tau} \int_0^\tau ds e^{-iH_0 s} (V - \prod_{H_0} V) e^{iH_0 s} = - \sum_{\nu, j, j', j''} \langle \nu, j | V | \nu, j' \rangle \langle \nu, j' | \langle \nu, j'' |,
\tag{5b}
\]

where \( \nu \) labels the different eigenvalues \( E^{(0)}_\nu \) of \( H_0 \), and \( j \) is a degeneracy index which distinguishes different basis vectors \( |\nu, j \rangle \) of the degeneracy eigenspace, and the units have been chosen such that \( h = 1 \). The operator \( \prod_{H_0} \) is the projector on the kernel of the application \( A \to [H_0, A] \). We can iterate the KAM procedure by the second transformation \( e^{iW_2} \) taking \( H_0 + \epsilon D_1 \) as the new reference Hamiltonian and \( \epsilon^2 V_2 \) as the new perturbation:

\[
e^{-\epsilon W_2} (H_0 + \epsilon D_1 + \epsilon^2 V_2) e^{\epsilon W_2} = H_0 + \epsilon D_1 + \epsilon^2 D_2 + \epsilon^4 V_3 \tag{6}
\]

where \( [D_2, H_0] = [D_2, D_1] = 0 \). Notice that the new perturbation is not of order as \( \epsilon^3 \) as would be the case in a standard perturbation theory, but is of order \( \epsilon^4 \). Similarly after \( n \) iteration we obtain:

\[
e^{-\epsilon W_n} \cdots e^{-\epsilon W_2} (H_0 + \epsilon D_1 + \epsilon^2 V_2) e^{\epsilon W_2} \cdots e^{-\epsilon W_1} = H_0 + \epsilon D_1 + \epsilon^2 D_2 + \cdots + \epsilon^{2n-1} D_n + \epsilon^{2n} V_{n+1}, \tag{7}
\]

where \( D_n, W_n, \) and \( V_{n+1} \) are calculated in a similar way as eqs. (5a,b) and(3 ) by

\[
D_n = \prod_{H_0^{n-1}} V_n,
\]

\[
W_n = \lim_{x \to +\infty} \frac{1}{x} \int_0^x ds e^{-iH_0'^{n-1} s} (V_n - \prod_{H_0} V_n) e^{iH_0'^{n-1} s},
\]

\[
\epsilon^{2n} V_{2n+1} = \sum_{m=2}^{\infty} \frac{\epsilon^{m2n-1}}{m!} ((m-1)L_m^{m-1} W_{n+1} + \frac{\epsilon^{m-1}}{m!} D_n), \tag{8}
\]

With

\[
H_0^{n-1} = H_0 + \epsilon D_1 + \epsilon^2 D_2 + \cdots + \epsilon^{2n-2} D_{n-1}. \tag{9}
\]

Since the perturbation term after \( n \)-th KAM iteration is order of \( \epsilon^{2n} \) (in contrast with the order of \( \epsilon^n \) in the standard perturbation theory), the KAM algorithm is called superconvergent. In the following discussion, we do not write explicitly the ordering parameter \( \epsilon \).

A resonance is defined as a degeneracy of an eigenvalue \( E^{(0)}_\nu \) of \( H_0 \) and is said to be active if the perturbation \( V \) has nonzero matrix elements in the degeneracy subspace of \( E^{(0)}_\nu \), \( \langle \nu, j | V | \nu, j' \rangle \neq 0 \) for some \( j, j' \). Otherwise the resonance is called passive or mute. In the case of quasi-resonance, the perturbative expansion cannot be expected to converge [9,10]. Indeed the perturbative formulas diverge close to resonances due to the appearance of small denominators. Using the method presented in Ref. [10], such divergences can be avoided.

The concept of resonance is defined intrinsically with respect to \( H_0 \), while the distinction between active and passive depends on the relation between \( H_0 \) and \( V \). The analysis of resonances involves thus three aspects: (i) Decomposition of the Hamiltonian into \( H = H_0 + V \). (ii) Determination of degenerate eigenvalues of the unperturbed Hamiltonian. (iii) Detection of the resonant terms in the perturbation \( V \) that couple these degenerate eigenstates.

The resonance terms of \( V \) can be detected by projectors of type \( \prod_{H_0} \) that extract a block-diagonal part of \( V \) relative to \( H_0 \), where the blocks are generated by the degeneracy subspaces. In the absence of active resonance, when all the eigenvalues of \( H_0 \) are non-degenerate or when the resonances are mute, the matrix
representation of $\Pi H_0$ is in fact diagonal in the
eigenbasis of $H_0$. In presence of active resonances, the
block-diagonal effective Hamiltonian that take into
account the considered resonance of the original
Hamiltonian can be written as

$$H^{\text{eff}} = H_0 + \Pi H_0 V. \quad (10)$$

We will call the transformation that diagonalizes $H^{\text{eff}}$
$\text{Resonant Transformation}$ (RT). The Hamiltonian is
transformed under $\mathcal{R}$ as follows:

$$H_1 = \mathcal{R}^{\dagger} H \mathcal{R} = \mathcal{R}^{\dagger} H^{\text{eff}} \mathcal{R} + \mathcal{R}^{\dagger} (V - \Pi H_0 V) \mathcal{R}$$

$$= H_1^{(0)} + V_1, \quad (11)$$

Where $H_1^{(0)}$ is defined as the new renormalized
reference Hamiltonian and $V_1$ is the new perturbation.
The effect of $\Pi H_0 V$ in eq. (10) is to lift the degeneracy of
$H_0$ (completely or partially). This can happen in two
ways: either the active resonance is transformed into a
passive one (e.g. in the case of zero-field resonances) or
the resonance disappears completely (when a crossing
is transformed into an avoided crossing). The new
Hamiltonian $H_1$ can however have other resonance at
different values of the coupling parameter. If $H_1^{(0)} + V_1$
do not have any other active resonance in the
considered range of the coupling constant, we can at a
second stage improve the spectrum by a perturbative
expansion which is expected to converge. If there are
active resonances, we have to iterate the renormalization
procedure by applying another RT.

3. Semiclassical Floquet formalism

In this section we overview the Floquet formalism
[3,10,15] in order to provide an essential tool which will
be used in the next sections. We consider the atom laser
interaction in a dipole coupling approximation. When the
time dependence of the periodic Hamiltonian is
introduced through the time evolution of the initial
phase, i.e. the amplitude and frequency of the laser field
are time-independent, the Floquet formalism is just
mathematically convenient tool that allows us to
transform the Schrödinger equation with a time-
dependent Hamiltonian into an equivalent equation with
a time-independent Hamiltonian. This new equation is
defined on an enlarged Hilbert space. The time
dependence has been substituted by the introduction of
one auxiliary dynamical variable for each laser
frequency. The semiclassical Hamiltonian can be written as

$$H = H(\theta_0 + \omega t) = H_0 - \mu \varepsilon \cos(\theta_0 + \omega t), \quad (12)$$

where $\mu$ is the dipole moment and $H_0$ the Hamiltonian
of the free molecule, $\varepsilon$ is the amplitude of the electric
field and $\omega$ its frequency. The corresponding
Schrödinger equation

$$i\hbar \frac{\partial}{\partial t} \Phi = H(\theta_0 + \omega t) \Phi, \quad \Phi \in \mathcal{H} \quad (13)$$

is defined on a Hilbert space $\mathcal{H}$, which can be of infinite
dimension (e.g. the space of square-integrable function
$\mathcal{H} = L^2(\mathbb{R}^3, d^3x)$, where $n$ is the number of the degrees of
freedom of the molecule) or infinite dimension (e.g. in
N-level models $\mathcal{H} = \mathbb{C}^N$). In eq. (13) the initial phase $\theta_0$
appears as a parameter. One can think of eq. (13) as a
family of equation parameterized by the angle $\theta_0$. We
denote the corresponding family of propagators by
$U(t,t_0;\theta_0)$, which describe the time evolution of
arbitrary initial condition $\Phi(t_0)$:

$$\Phi(t) = U(t,t_0;\theta_0)\Phi(t_0), \quad (14)$$

and satisfy

$$i\hbar \frac{\partial}{\partial t} U(t,t_0;\theta_0) = H(\theta_0 + \omega t)U(t,t_0;\theta_0),$$

$$U(t,t;\theta_0) = 1_{\mathcal{H}}. \quad (15)$$

The Floquet Hamiltonian $K$, also called $\text{quasienergy}$
operator, is constructed as follows: we define an
enlarged Hilbert space

$$\mathcal{K} = \mathcal{H} \otimes \mathcal{L}, \quad (16)$$

where $\mathcal{L} = L^2(\mathbb{S}^1, d\theta/2\pi)$ denotes the space of square
integrable function on the circle $\mathbb{S}^1$ of length $2\pi$, with a scalar
product

$$\langle f_1 | f_2 \rangle_{\mathcal{L}} = \int_{\mathbb{S}^1} \frac{d\theta}{2\pi} f_1^*(\theta) f_2(\theta). \quad (17)$$

This space is generated by the orthonormal basis
$$\{ e^{i k \theta}, k \in \mathbb{Z} \}.$$

We first lift the family of operators $U(t,t_0;\theta_0)$
(denoted on $\mathcal{H}$) into an operator $U(t,t_0)$ acting on the
enlarged space $\mathcal{K}$ by considering the initial phase $\theta_0$
as a multiplication operator $\theta$. This operator is unitary in
$\mathcal{K}$. On the enlarged Hilbert space $\mathcal{K}$ the Floquet
Hamiltonian $K$ is defined as:

$$K = -i\hbar \omega \frac{\partial}{\partial \theta} + H. \quad (18)$$

In this expression $H$ is just the semiclassical Hamiltonian
(12) but with the phase taken as a dynamical variable.
The usefulness of Floquet Hamiltonian comes from the
fact that is time-independent and that the dynamics it
defines on $\mathcal{K}$ is essentially equivalent with the one of eq.
(13). The Floquet Hamiltonian $K$ defines a time
evolution in $\mathcal{K}$ through the equation

$$i\hbar \frac{\partial}{\partial t} \Psi = K \Psi, \quad \Psi \in \mathcal{K} = \mathcal{H} \otimes \mathcal{L}. \quad (19)$$

This time evolution can be expressed in terms of propagator $U_{\mathcal{K}}$
characterized by
\[ i\hbar \frac{\delta}{\partial t} U_K(t, t_0) = KU_K(t, t_0), \quad U_K(t, t) = 1_K, \] 
\[ i\hbar \frac{\delta}{\partial t} U_K(t, t_0) = KU_K(t, t_0), \quad U_K(t, t) = 1_K. \]
i.e. \( \Psi(t) = U_K(t, t_0) \Psi(t_0) \), (where \( 1_K \) is the identity operator in \( K \)). Since \( K \) is time-independent, the propagator can be written as
\[ U_K(t, t_0) = U_K(t - t_0) = e^{-iK(t-t_0)/\hbar}. \]
In order to establish a relation between \( U \) and \( U_K \), we define the following phase translation operator \( T_{\alpha} \), which acts on \( f \in \mathcal{L} \), by
\[ T_{\alpha} f(\theta) := f(\theta + \alpha \theta), \]
and can be expressed as
\[ T_{\alpha} = e^{i\alpha \theta \partial / \partial \theta}. \]
The relation between \( U \) and \( U_K \) can be expressed by
\[ T_{-\alpha} U T_{\alpha} = U_K(t - t_0) = e^{-iK(t-t_0)/\hbar}. \]
It implies that if is a solution of eq. (19) then we can obtain a solution of eq. (13) by
\[ \Phi(t) = [T_{\alpha} \Psi(t)]_{\theta = 0} = \Psi(t, \theta + \alpha \theta), \]
where fixing a constant value for the dynamical variable means returning into the original Hilbert space \( \mathcal{H} \).

4. Two-level atom in a one-mode classical field
In this section, we consider the model of a two-level atom coupled with a single-frequency laser field. The Floquet Hamiltonian acting on the enlarged Hilbert space \( \mathcal{K} = \mathcal{L} \otimes \mathcal{H}, \mathcal{H} = \mathbb{C}^2 \), is given by
\[ K = -i\alpha \frac{\partial}{\partial \theta} \otimes 1_2 + 1_2 \otimes \frac{\partial \theta}{2} \sigma_z + \Omega \cos \theta \otimes \sigma_y, \]
where \( \Omega = -\mu \omega \) is the coupling of the atomic transition coupled by laser field, and \( \sigma_z, \sigma_y \) are Pauli matrices:
\[ \sigma_z = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \sigma_y = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \sigma_x = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \]
where \( 1_2 \) is the 2 x 2 identity matrix. This Hamiltonian is identical with the Hamiltonian of a two-level atom in a one-mode quantized field,
\[ H = \omega a^\dagger a \otimes 1_2 + 1_2 \otimes \frac{\partial \theta}{2} \sigma_z + g(a + a^\dagger) \otimes \sigma_x, \]
in the limit of infinite cavity volume, large photon number average, and constant photon density [15]. Since \( \cos \theta = (e^{-i\theta} + e^{i\theta}) / 2 \), comparing eqs. (26) and (28), we can establish a correspondence among the operators \( a, a^\dagger \) and \( e^{i\theta}, e^{-i\theta} \). This point will be clarified further in the next section.

The parity operator of this system is
\[ P = e^{+i\frac{\pi}{2} \theta \partial / \partial \theta} \otimes \sigma_z, \]
with the properties
\[ [\sigma_z, \sigma_y] = 0, \quad P = P^\dagger, \quad P^2 = 1_2 \otimes 1_2. \]
As a consequence, the eigenstate of \( K \) can be separated into symmetry classes, even or odd, under \( P \):
\[ P | \phi_k, \pm \rangle = \pm | \phi_k, \pm \rangle, \quad K | \phi_k, \pm \rangle = E_{k, \pm} | \phi_k, \pm \rangle. \]
The conceptual framework for the solution of the solution of this system based on the construction of unitary transformation can be described as follows: First, we decompose the Hamiltonian in two terms as \( H = H_0 + V \). Depending on the considered range of the parameters of the system, different decompositions may be considered. \( H_0 \) is an \textit{a priori} operator that is a regular function exclusively of the operators \( N = -i \alpha \partial / \partial \theta \) (photon-number operator) and \( \sigma_z \). The operators \( N \) and \( \sigma_z \) can be considered in the present model as quantum analogues of classical global actions [16], and \( H_0 \) can be labelled \textit{integrable}. The perturbation \( V \) contains functions that involve also the other operators \( e^{-i\theta}, e^{i\theta}, \sigma_x, \sigma_y \). The goal is to determine a unitary transformation \( U \), that should be expressed in terms of well-behaved regular functions of \( e^{-i\theta}, e^{i\theta}, \sigma_x, \sigma_y, \sigma_z \), such that:
\[ U^\dagger (H_0(N, \sigma_z) + V(e^{-i\theta}, e^{i\theta}, \sigma_x, \sigma_y, \sigma_z)) U = H'(N, \sigma_z), \]
where \( H' \) is a regular function \( f \) exclusively of the action operators \( N, \sigma_z : H'(N, \sigma_z) = f(n, \sigma_z) \). With this transformation the eigenvectors of \( H \) can be expressed as \( | \phi_{k, \pm} \rangle = U(k) \otimes | \pm \rangle \) and the corresponding eigenvalues as \( E_{k, \pm} = f(k, \pm 1) \) where \( N | k \rangle = k | k \rangle \) and \( \sigma_z | \pm \rangle = \pm | \pm \rangle \).

Most of the perturbative approaches can be interpreted as methods to find approximation of the transformation of \( U \). The presence of resonance is one of the central difficulties in the construction of \( U \), as will be made precise below. In this paper we discuss an iterative approach that consist of constructing first some approximation of \( U \) that take into account the dominating effects of certain number of resonances. The transformations involved in this stage are far from the identity and have a clearly non-perturbative character. Once we have a transformation that takes into account the main effect of a set of resonances that are relevant in a considered interval of coupling constant \( \Omega \), a perturbative approach can be applied to improve the approximation quantitatively. The transformations involved in this second stage can be considered as deformations of the identity, since they can be written in the form \( e^{iW} \). This stage cannot be implemented if the resonances are not taken care of beforehand.
As in classical mechanics, the construction $U$ leading to Hamiltonian contains only action variables can often be considered in two steps: $U = U_1 U_2$. In the first step, that is called reduction, the Hamiltonian is transformed by into a form that contains functions of $\sigma_z, \sigma_x, \sigma_y$ and $N$, but not of $e^{-i\theta}$ and $e^{i\theta}$. The degree of freedom of the field is made trivial and the number of non-trivial degrees of freedom is thus reduced by one. When we apply this reduction to the effective Hamiltonian, we obtain a reduced effective Hamiltonian. We remark that in the literature, this reduced effective Hamiltonian is often called simply effective Hamiltonian.

In the second step, the reduced Hamiltonian is transformed under $U_2$ into a form that contains functions of only $N$ and $\sigma_z$. For the model (26), the reduction step corresponds to a diagonalization in the space $\mathcal{L}$ and the second step corresponds to a diagonalization in the atomic Hilbert space which in this case trivial. The construction of the RT is based on this reduction procedure.

5. Effective Hamiltonians in the weak - coupling regime

The Floquet Hamiltonian (26) in the weak-coupling regime can be decomposed as follows:

$$K_0 = -i\omega \frac{\partial}{\partial \theta} + \frac{\omega_0}{2} \sigma_z, \quad V = \Omega \cos \theta \sigma_z.$$  

(33)

The eigenvalues and eigenvectors of $K_0$ are

$$E_k^{(0)} = k\omega + \omega_0/2,$$

$$\phi^{(0)}_{k,\pm} = |k, \pm\rangle, \quad \langle k, + | = 0 \begin{pmatrix} k \rangle \langle 0 \end{pmatrix}, \quad \langle k, - | = \begin{pmatrix} 0 \rangle \langle k \end{pmatrix},$$

(34)

where $|k\rangle$ is eigenvector of the relative photon number operator $-i\frac{\partial}{\partial \theta}$ as $-i\frac{\partial}{\partial \theta}|k\rangle = k|k\rangle$ with orthonormal condition

$$\langle k | k \rangle = \int_0^{2\pi} \frac{d\theta}{2\pi} e^{-i(k-k')\theta} = \delta_{k,k'}.$$  

(35)

For the case of one-photon resonance $\omega = \omega_0$, the eigenvalues of $K_0$ are degenerate as $E_k^{(0)} = E_{k+1}^{(0)}$. The degeneracy eigenspaces are spanned by vectors $\phi_{k+,0}$ and $\phi_{k+1,-}$. The resonant part of $V$ is obtained by eq. (5a):

$$V_{res} = \Pi_{K_0} V = \sum_{k=-\infty}^{\infty} \langle k, + | V | k+1, - \rangle \langle k+1, - | + \langle K+1, - | V | k, + \rangle \langle k, + |$$

$$= \Omega \frac{1}{2} \begin{pmatrix} 0 & e^{-i\theta} \\ e^{i\theta} & 0 \end{pmatrix},$$

(36)

where we have obtained the resonant elements of $V$ as follows:

$$\langle k+1, - | V | k, + \rangle = \Omega \langle k+1 | \cos \theta | k \rangle \langle - \sigma_x |$$

$$= \Omega \langle k+1 | \cos \theta | k \rangle$$

$$= \Omega \int_0^{2\pi} \frac{d\theta}{2\pi} e^{-i(k+1)\theta} \cos \theta e^{i\theta}$$

$$= \Omega/2,$$

(37)

and

$$\sum_k |k+1, - \rangle \langle k+1, - | = e^{i\theta} \begin{pmatrix} 0 & 0 \\ 1 & 0 \end{pmatrix}. $$

(38)

The spectral representation of the multiplication operator $e^{i\theta}$ and the creation operator $a^+$ can be compared as follows:

$$e^{i\theta} = \sum_{n=-\infty}^{\infty} |k+1\rangle \langle k| \leftrightarrow a^+ = \sum_{n=-\infty}^{\infty} \sqrt{n+1} |n+1\rangle \langle n|$$

(39)

Therefore, the effective Floquet Hamiltonian containing the one-photon resonance (corresponding to the Jaynes-Cummings Hamiltonian [1]) is

$$K_0^{eff} = -i\omega \frac{\partial}{\partial \theta} + \frac{\omega}{2} \sigma_z + \frac{\Omega}{2} \begin{pmatrix} 0 & e^{-i\theta} \\ e^{i\theta} & 0 \end{pmatrix}.$$  

(40)

Next we transform $K_0^{eff}$ by a RT to a regular function of exclusively the action operators $-i\omega \frac{\partial}{\partial \theta} \sigma_z$. The reduction step of RT denoted $R$ is a unitary (as opposed to isometric in the quantized field [9]) transformation that can eliminate the $\theta$- dependence in the resonant term,

$$R V_{res} R = \frac{\Omega}{2} \sigma_x.$$  

(41)

The suitable transformation satisfying this condition is

$$R = \begin{pmatrix} e^{i\theta} & 0 \\ 0 & 1 \end{pmatrix} = \sum_{k=-\infty}^{\infty} \begin{pmatrix} k \rangle \langle k+1 | \langle k+1, - | \end{pmatrix}.$$  

(42)

Applying $R$ on $K$ gives:

$$R^\dagger K R = -i\omega \frac{\partial}{\partial \theta} + \frac{\omega}{2} \sigma_z + \frac{\Omega}{2} \begin{pmatrix} 0 & e^{2i\theta} \\ e^{-2i\theta} & 0 \end{pmatrix}$$

(43)

where we have used the relation

$$\frac{\partial}{\partial \theta} R = R \frac{\partial}{\partial \theta} + \frac{\partial R}{\partial \theta}.$$  

(44)

We remark that when there is a small detuning, $\omega_0 = \omega + \Delta$, we can decompose $K$ as

$$K_0 = -i\omega \frac{\partial}{\partial \theta} + \frac{\omega}{2} \sigma_z \quad \text{and} \quad V = \frac{\Delta}{2} \sigma_z + \Omega \cos \theta \sigma_x, \quad \text{where}$$

the detuning has been included in the perturbation. This decomposition leads to the same resonant term as eq. (36), and thus the eq. (43) is modified by adding the term
\[ \Delta \sigma_z. \] The corresponding reduced effective Hamiltonian in the atomic Hilbert space and in the rotating-wave approximation is
\[ H_{\text{eff}}^{\text{rot}} = \begin{pmatrix} \Delta & \Omega \\ \Omega & -\Delta \end{pmatrix}. \] (45)
We note that the Rabi frequency of a system having a reduced effective Hamiltonian such as eq. (45) is
\[ r = \Omega / 2. \] (46)
This means that if the system starts from the state \( \mid + \rangle \) or \( \mid - \rangle \), the system will populate periodically at later times the states \( \mid + \rangle \) and \( \mid - \rangle \) with period of \( 2\pi / \omega \). On the other hand, we will see in section VI that the Rabi frequency of this system in the strong-coupling regime is different from eq. (46).
The second step of the first RT consists in diagonalizing the \( \sigma_x \) term in \( R^\dagger K_0^{\text{eff}} R \) by a \( \pi / 2 \) rotation around the \( y \)-axis,
\[ T = e^{\frac{i}{4} \sigma_y^2} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & -1 \\ 1 & 1 \end{pmatrix}, \] (47)
with the properties
\[ R^\dagger \sigma_z R = \sigma_z, \quad R^\dagger \sigma_x R = -\sigma_x, \] (48)
as follows:
\[ K_1 = R^\dagger K_0^{\text{eff}} R = -i\omega \frac{\partial}{\partial \theta} \frac{\omega}{2} \sigma_x + \frac{\Omega}{2} \begin{pmatrix} \cos 2\theta & i\sin 2\theta \\ -i\sin 2\theta & \cos 2\theta \end{pmatrix}. \] (49)
Now, we can decompose the Hamiltonian \( K_1 \) to:
\[ K_0 = -i\omega \frac{\partial}{\partial \theta} \frac{\omega}{2} \sigma_x, \quad V_1 = \frac{\Omega}{2} \begin{pmatrix} \cos 2\theta & i\sin 2\theta \\ -i\sin 2\theta & \cos 2\theta \end{pmatrix}, \] (50)
and the resonant term of \( V_1 \) is extracted by \( \prod_{K_1^{(0)}} V_1 \) in terms of the degeneracies of \( K_1^{(0)} \). The resonances of \( K_1 \) are nonlinear as the degeneracies of \( K_1^{(0)} \) are \( \Omega \)-dependent. This procedure can be iterated to obtain effective Hamiltonians for larger values of the coupling constant \( \Omega \) [9,10].

6. Effective Hamiltonians in the strong - coupling regime
In this section we consider the Floquet Hamiltonian (26) in the strong-coupling regime \( \Omega >> \omega_0 \) and we extend the ideas given in Refs. [9, 10] to the case of multiphoton resonances in Floquet formalism. The strong-coupling regime suggests an alternative decomposition as
\[ K_0 = -i\omega \frac{\partial}{\partial \theta} + \Omega \cos \theta \sigma_x, \quad V = \frac{\omega_0}{2} \sigma_z \] (51)
where \( K_0 \) contains again all the unbounded operators of the system, and the perturbation \( V \) is a bounded operator. We take an approach which is very similar to the one of Ref. [9]. In the case of one-photon resonance \( \omega_0 = \omega \), this approach allows us to obtain an effective Hamiltonian which describes the essential structure of the system for the whole range of values of the coupling \( \Omega \in [0, \infty) \). However for multi-photon resonances \( \omega_0 = m\omega, \ m = 3, 5, 7, \ldots \), this approach works only for large values of \( \Omega \). As in [9], we first apply the transformation (47) to diagonalize \( K_0 \) in the atomic Hilbert space
\[ T^\dagger KT = -i\omega \frac{\partial}{\partial \theta} + \Omega \cos \theta \sigma_z - \frac{\omega_0}{2} \sigma_x. \] (52)
Next we apply a second unitary transformation to eliminate the \( \theta \)-dependence in \( T^\dagger K_0 T \):
\[ U = \begin{pmatrix} \Omega \sin \theta & 0 \\ 0 & e^{i\omega \sin \theta} \end{pmatrix}, \] (53)
which leads to
\[ U^\dagger T^\dagger K T U = -i\omega \frac{\partial}{\partial \theta} - \frac{\omega_0}{2} \begin{pmatrix} 0 & e^{i\gamma \sin \theta} \\ e^{-i\gamma \sin \theta} & 0 \end{pmatrix}, \] (54)
where
\[ \gamma = \frac{2\Omega}{\omega}. \] (55)
Now by using the relation
\[ e^{i\gamma \sin \theta} = \cos(\gamma \sin \theta) \pm i\sin(\gamma \sin \theta) \]
\[ = J_0(\gamma) + 2 \sum_{p=2,4,\ldots} J_p(\gamma) \cos p\theta \pm 2i \sum_{q=1,3,\ldots} J_q(\gamma) \sin q\theta \] (56)
we can write \( U^\dagger T^\dagger K T U \) in terms of harmonics of the new perturbation as
\[ U^\dagger T^\dagger K T U = -i\omega \frac{\partial}{\partial \theta} - \frac{\omega_0}{2} J_0(\gamma) \sigma_x - \frac{\omega_0}{2} \sum_{p=2,4,\ldots} J_p(\gamma) \cos p\theta \sigma_x + \omega_0 \sum_{q=1,3,\ldots} J_q(\gamma) \sin q\theta \sigma_y. \] (57)
To diagonalize the \( \sigma_x \) terms, we apply the transformation \( T \) again.
\[ K_1 = U^\dagger T^\dagger K T U = -i\omega \frac{\partial}{\partial \theta} - \frac{\omega_0}{2} J_0(\gamma) \sigma_z - \omega_0 \sum_{p=2,4,\ldots} J_p(\gamma) \cos p\theta \sigma_x + \omega_0 \sum_{q=1,3,\ldots} J_q(\gamma) \sin q\theta \sigma_y. \] (58)
At a second stage, we decompose \( K_1 = K_1^{(0)} + V_1 \) as
Figure 1. Comparison of the exact numerical eigenvalues (solid lines) of eq. (26) as a function of the coupling constant in the resonant case $\omega_0 = \omega$, with the approximate eigenvalues (dashed and dotted lines) obtained from eq. (69).

$$K^{(0)}_1 = -i\omega \frac{\partial}{\partial \theta} - \frac{\omega_0}{2} J_0(\gamma) \sigma_Z ,$$

$$V_1 = -\omega_0 \sum_{p=2,A} J_p(\gamma) \cos p \theta \sigma_Z + \omega_0 \sum_{q=3,\gamma} J_q(\gamma) \sin q \theta \sigma_y .$$

(59)

For

$$\omega_0 J_0(\gamma) = m \omega ,$$

(60)

The eigenvalues of $K^{(0)}_1$ are with degeneracies $\lambda_{k,+} = \lambda_{k-1,-}$. The eq. (60) is the nonlinear resonance condition as it depends on the coupling $\Omega$. For $\Omega \to 0$ this condition is reduced to the $m$-photon zero-field resonance condition $\omega_0 = m \omega$. The resonant part of $V_1$ can be calculated as follows:

$$\Pi_{K^{(0)}} V_1 = \sum_{k=-\infty}^{\infty} |k, m-1\rangle \langle k, m-1| V |k, +\rangle \langle k, +|$$

$$+ |k, +\rangle \langle k, +| V |k, m-1\rangle \langle k, m-1|$$

$$= -\frac{\omega_0}{2} J_m(\gamma) \begin{pmatrix} 0 & e^{i m \theta} \\ e^{-i m \theta} & 0 \end{pmatrix} , \quad m = 1, 3, 5, \ldots$$

(61)

where we have used the relations

$$\sum_{k=-\infty}^{\infty} |k, m-1\rangle \langle k, m-1| = e^{-i m \theta}$$

$$\sum_{k=-\infty}^{\infty} \langle k, m-1| k, m-1\rangle = e^{i m \theta} .$$

(62)

We observe that the terms with even $m$ are absent in the average of $V_1$ with respect to $K^{(0)}_1$. So the effective Floquet Hamiltonian of the system in strong-coupling regime can be written as

$$K^{eff}_1 = -i\omega \frac{\partial}{\partial \theta} + \frac{\omega_0}{2} J_0(\gamma) \sigma_z - \frac{\omega_0}{2} J_m(\gamma) \begin{pmatrix} 0 & e^{i m \theta} \\ e^{-i m \theta} & 0 \end{pmatrix} , \quad m = 1, 3, 5, \ldots$$

(63)

Figure 2. Comparison of the exact numerical eigenvalues (solid lines) of eq. (26) as a function of the coupling constant in the resonant case $\omega_0 = 3 \omega$, with the approximate eigenvalues (dashed and dotted lines) obtained from eq. (69).

The reduction step of the m-photon RT to eliminate the $\theta$-dependence of the resonant term in $K^{eff}_1$ can be performed by

$$R_m = \begin{pmatrix} 1 & 0 \\ 0 & e^{-i m \theta} \end{pmatrix} ,$$

(64)

which leads to

$$R_m^* K^{eff}_1 R_m = -i\omega \frac{\partial}{\partial \theta} + \begin{pmatrix} 0 & 0 \\ 0 & -m \omega \end{pmatrix} - \frac{\omega_0}{2} J_0(\gamma) \sigma_z - \frac{\omega_0}{2} J_m(\gamma) \sigma_x .$$

(65)

This gives the reduced effective Hamiltonian of the system, in the atomic Hilbert space, for strong-coupling regime:

$$H^{eff}_S = \begin{pmatrix} -\frac{\omega_0}{2} J_0(\gamma) & -\frac{\omega_0}{2} J_m(\gamma) \\ -\frac{\omega_0}{2} J_m(\gamma) & -m \omega + \frac{\omega_0}{2} J_0(\gamma) \end{pmatrix} .$$

(66)

Compare this Hamiltonian with eq. (45). In this case the Rabi frequency of the system, instead of $\Omega / 2$, is

$$r = \frac{\omega_0}{2} J_m(\frac{2 \Omega}{\omega}) .$$

(67)

One can see that for $\Omega \to 0$ and $m = 1$, the Hamiltonian (45) with zero detuning can be recovered as

$$J_0(\gamma) \to 1 , \quad J_m(\gamma) \to \left(\frac{\gamma}{2}\right)^m , \quad \text{for } \gamma \to 0 ,$$

(68)

The eigenvalues of eq. (66) can be calculated easily as

$$E_m = \frac{-m \omega}{2} \pm \frac{1}{2} \sqrt{m \omega - \omega_0 J_0(\gamma)^2 + [\omega_0 J_m(\gamma)]^2} ,$$

(69)

$$m = 1, 3, 5, \ldots$$

Figures 1, 2 and 3 compare the exact eigenvalues of the
system with the approximate eigenvalues obtained from eq. (69). We observe in figure 1 that for one-photon resonance the approximate eigenvalues are valid for the whole range of the coupling $\Omega$, while for the three- and five-photon resonances in figures 2, 3 the eigenvalues obtained from eq. (69) are valid only for large values of $\Omega$.

7. Conclusions and discussions

We have used a non-perturbative method based on the quantum averaging technique to determine the spectral properties of a semiclassical system (a two-level atom in a one-mode classical field) in two regimes of weak- and strong-coupling. It consists in the construction of unitary transformations that leads to an effective reduced Hamiltonian. These transformations are adapted to the structure of the resonances. Their role is to construct a first effective Hamiltonian that contains the main qualitative features of the spectrum in a given range of the coupling parameter. The diagonalized form of this effective Hamiltonian, which depends parametrically on the coupling constant, is then taken as a new reference Hamiltonian around which one can apply perturbative techniques to improve the quantitative accuracy of the spectrum.

We have analyzed the resonances in two regimes of weak and strong coupling. The analysis of the strong-coupling regime of this model leads to results that are valid for all values of the coupling and for all energies in the case of a one-photon resonance, $\omega_0 = \omega$ when this later resonance is additionally treated. The possibility to obtain such a global result is due to a particular property of the model, and one cannot expect to obtain it for general models. The particular property is that the part we selected as the reference Hamiltonian $H_0$ in the strong-coupling regime contains all the unbounded operators of the complete model and is explicitly solvable. The term that was left to be treated by RT and perturbation theory is a bounded operator.

The extension of the strong-coupling regime for multi-photon resonances $\omega_0 = m\omega$, $m = 3, 5, 7 \ldots$ to the weak-coupling regime is under consideration. This can be related to the fact that in the weak-coupling decomposition of the complete Hamiltonian (26), the multi-photon resonances are passive. These resonances could be treated after applying a KAM-type transformation [10, 17] to transform them from passive to active resonances.

References

18. The parameter $\varepsilon$ can be thought as a small parameter that can be used to attribute an order of perturbation to the terms appearing in each iteration of the KAM-type perturbation theory.