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Introduction

This volume contains a selection of invited and contributed research papers presented at the VIII International Readings on Quantum Optics (IRQO'99) held in Kazan (which is capital of Tatarstan Republic, Russian Federation) October 27-29, 1999. Such International Readings are traditional and annual. The main aim of IRQO'99 was the discussion of the modern art of investigations on the following broad topics: coherent states of electromagnetic field; quantum correlations and photons' statistics; squeezed states of light; quantum effects in nonlinear optics; laser cooling of gases and solids; optical superradiance and other collective processes; photon echo and other transient phenomena; optical phase memory, based on these phenomena; problems of gamma-optics and gamma-lasers and also was analysis of the ways of the scientifical and practical use of the optical coherent and quantum phenomena.

Now the quantum optics is the perspective direction of modern physics. The vacuum fluctuations, quantum beats, coherent and squeezed states of light, field-field and photon-photon interferometry, photon antibunching, counting and photon statistics, lasing without inversion, coherent trapping, correlated spontaneous emission laser, holographic laser, the two-photon correlated emission laser, the Einstein-Podolsky-Rosen paradox, teleportation, quantum eraser, quantum optical Ramsey fringes, micromaser, laser cooling, superradiance and optical transient phenomena are the most interesting scientifical problems of quantum optics. All these problems were discussed on IRQO'99. Total number of papers was equal to 88, including 28 invited papers. About 100 scientists from eight countries (Russia, Belarus, Ukraine, Japan, Belgium, France, Switzerland, Greece) presented their recent achievements at IRQO'99.

The IRQO'99 Program included the papers, devoted to peculiarities of detection of the slow coherent optical signals, "hidden" light polarization, chaos and squeezing in quantum optics, quantum teleportation, nonlocal interactions and quantum dynamics, quantum properties of system "atom plus cavity", interaction of atom with strong laser fields. The problems of the optical Dicke superradiance, photon echoes, self-induced transparency and other coherent phenomena were actively discussed also. The work of the gamma-optics section evoked the essential interest of participants and, particularly, the papers of Prof. T.Arisawa (Japan) and Prof. Jos Odeurs with collaborators (Belgium).

The Proceedings of IRQO'99 will be of interest to a broad spectrum of the international technical community since the area of quantum optics is modern and it has prospective scientific and practical use.

Vitaly V.Samartsev

Contents

SESSION 1

Problems of Quantum Optics

Collective Phenomema in the Interaction of Radiation with Matter

V.I. Yukalov

Insituto de Fisica de São Carlos, Universidade de São Paulo
Caixa Postal 369, São Carlos, São Paulo 13560–970, Brazil

Abstract

The aim of this communication is to present in a concentrated form the main ideas of a method, developed by the author, for treating strongly nonequilibrium collective phenomena typical of the interaction of radiation with matter, as well as to give a survey of several applications of the method. The latter is called the Scale Separation Approach since its basic techniques rely on the possibility of separating different space–time scales in nonequilibrium statistical systems. This approach is rather general and can be applied to diverse physical problems, several of which are discussed here. These problems are: Superradiance of nuclear spins, filamentation in resonant media, semiconfinement of neutral atoms, negative electric current, and collective liberation of light.

1 Introduction

Strongly nonequilibrium processes that occur in statistical systems and involve their interaction with radiation are usually described by complicated nonlinear differential and integro–differential equations [1–3]. For treating these difficult problems, a novel approach has recently been developed [4–7] called the *scale separation approach* since its main idea is to formulate the evolution equations in such a form where it could be possible to separate several characteristic space–time scales. In many cases, different scales appear rather naturally being directly related to the physical properties of the considered system.

The scale separation approach has been employed for solving several interesting physical problems related to strongly nonequilibrium processes occurring under the interaction of radiation with matter. As an illustration, the following phenomena are selected for this report: *Superradiance of Nuclear Spins*, *Filamentation in Resonant Media*, *Semiconfinement of Neutral Atoms*, *Negative Electric Current*, and *Collective Liberation of Light*.

Since the scale separation approach makes the mathematical foundation for the following applications, its general scheme is described in Section 2. In Sections 3 to 7 concrete physical effects are briefly reviewed and the most important results are summarized.

2 Scale Separation Approach

Because of the pivotal role of this approach for treating different physical problems, its general scheme will be presented here in an explicit way [4–7]. It is possible to separate the following main steps, or parts, of the approach.

2.1 Stochastic quantization of short–range correlations

When considering nonequilibrium processes in statistical systems, one needs to write evolution equations for some averages $< A_i >$ of operators $A_i(t)$ where t is time and $i = 1, 2, \ldots, N$ enumerates particles composing the considered system. For simplicity, a discrete index i is used, although everywhere below one could mean an operator $A(\vec{r}_i, t)$ depending on a continuous space variable $\vec{r}_i$.

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There is the well known problem in statistical mechanics consisting in the fact that writing an evolution equation for $< A_i >$ one does not get a closed system of equations but a hierarchical chain of equations connecting correlation functions of higher orders. Thus, an equation for $< A_i >$ contains the terms as $\sum_j < A_i B_j >$ with double correlators $< A_i B_j >$, and the evolution equations for the latter involve the terms with tripple correlators, and so on. The simplest way for making the system of equations closed is the mean-field type decoupling $< A_i B_j > \rightarrow < A_i >< B_j >$. When considering radiation processes, this decoupling is called the semiclassical approximation. Then the term $\sum_j < A_i B_j >$ reduces to $< A_i > \sum_j < B_j >$, so that one can say that $< A_i >$ is subject to the action of the mean field $\sum_j < B_j >$. The semiclassical approximation describes well coherent processes, when long-range correlations between atoms govern the evolution of the system, while short-range correlations, due to quantum fluctuations, are not important. However, the latter may become of great importance for some periods of time, for example, at the beginning of a nonequilibrium process when long-time correlations have had yet no time to develop. Then neglecting short-range correlations can lead to principally wrong results.

To include the influence of short-range correlations, the semiclassical approximation can be modified as follows:

$$\sum_j < A_i B_j > = < A_i > \left(\sum_j < B_j > + \xi \right) , \tag{1}$$

where ξ is a random variable describing local short-range correlations. It is natural to treat ξ as a Gaussian stochastic variable with the stochastic averages

$$\ll \xi \gg = 0 , \qquad \ll |\xi|^2 \gg = \sum_j | < B_j > |^2 , \tag{2}$$

where the second moment is defined so that to take into account incoherent local fluctuations. Since short-range correlations are often due to quantum fluctuations, the manner of taking them into account by introducing a stochastic variable ξ can be called the stochastic quantization. Then the decoupling (1) may be termed the *stochastic semiclassical approximation*. This kind of approximation has been used for taking into account quantum spontaneous emission of atoms in the problem of atomic superradiance.

2.2 Separation of solutions onto fast and slow

The usage of the stochastic semiclassical approximation makes it possible to write down a closed set of stochastic differential equations. The next step is to find such a change of variables which results in the possibility of separating the functional variables onto fast and slow, so that one comes to the set of equations having the form

$$\frac{du}{dt} = f(\varepsilon, u, s, \xi, t) , \qquad \frac{ds}{dt} = \varepsilon\, g(\varepsilon, u, s, \xi, t) , \tag{3}$$

where $\varepsilon \ll 1$ is a small parameter, such that

$$\lim_{\varepsilon \to 0} f \neq 0 , \qquad \lim_{\varepsilon \to 0} \varepsilon\, g = 0 . \tag{4}$$

As is evident, dealing with only two functions, u and s, and one small parameter ε is done just for simplicity. All procedure is straightforwardly applicable to the case of many functions and several small parameters.

From Eqs. (3) and (4) it follows that

$$\lim_{\varepsilon \to 0} \frac{du}{dt} \neq 0 , \qquad \lim_{\varepsilon \to 0} \frac{ds}{dt} = 0 , \tag{5}$$

which permits one to classify the solution u as fast, compared to the slow solution s. In turn, the slow solution s is a *quasi-invariant* with respect to the fast solution u.

The above classification of solutions onto fast and slow concerns time variations. In the case of partial differential equations, one has, in addition to time, a space variable $\vec{r}$. Then the notion of fast or slow functions can be generalized as follows [8,9]. Let $\vec{r} \in \mathbf{V}$, with $\text{mes}\mathbf{V} \equiv V$, and $t \in [0,T]$, where T can be infinite. Assume that

$$\lim_{\varepsilon\to 0} \ll \frac{1}{V}\int_{\mathbf{V}} \frac{\partial u}{\partial t}\, d\vec{r} \gg \neq 0\, , \qquad \lim_{\varepsilon\to 0} \ll \frac{1}{T}\int_0^T \vec{\nabla} u\, dt \gg \neq 0\, , \tag{6}$$

while

$$\lim_{\varepsilon\to 0} \ll \frac{1}{V}\int_{\mathbf{V}} \frac{\partial s}{\partial t}\, d\vec{r} \gg = 0\, , \qquad \lim_{\varepsilon\to 0} \ll \frac{1}{T}\int_0^T \vec{\nabla} s\, dt \gg = 0\, . \tag{7}$$

Then the solution u is called *fast on average*, with respect to both space and time, as compared to s that is *slow on average*. In such a case s is again a quasi–invariant as compared to u. In general, it may, of course, happen that one solution is fast with respect to time but slow in space, or vice versa, when compared to another function. The notion of quasi–invariants with respect to time is known in the Hamiltonian mechanics where they are also called adiabatic invariants. Here this notion is generalized to the case of both space and time variables [8,9].

2.3 Averaging method for multifrequency systems

After classifying in Eqs. (4) the function u as fast and s as slow, one can resort to the Krylov–Bogolubov averaging technique [10] extended to the case of multifrequency systems. This is done as follows.

Since the slow variable s is a quasi–invariant for the fast variable u, one considers the equation for the fast function u, with the slow one kept fixed,

$$\frac{\partial X}{\partial t} = f(\varepsilon, X, z, \xi, t)\, . \tag{8}$$

Here z is treated as a fixed parameter. The solution to Eq. (8), that is

$$X = X(\varepsilon, z, \xi, t)\, , \qquad z = const\, , \tag{9}$$

has to be substituted into the right–hand side of the equation for the slow function, and for this right–hand side one defines the average

$$\overline{g}(\varepsilon, z) \equiv \ll \frac{1}{\tau}\int_0^\tau g(\varepsilon, X(\varepsilon, z, \xi, t), z, \xi, t)\, dt \gg\, , \tag{10}$$

in which τ is the characteristic oscillation time of the fast function. In many cases, it is possible to take $\tau \to \infty$, especially when the period of fast oscillations is not well defined [2]. Then one comes to the equation

$$\frac{dz}{dt} = \varepsilon\, \overline{g}(\varepsilon, z) \tag{11}$$

defining a solution

$$z = z(\varepsilon, t)\, . \tag{12}$$

Substituting the latter into X, one gets

$$y(\varepsilon, \xi, t) = X(\varepsilon, z(\varepsilon, t), \xi, t)\, . \tag{13}$$

The pair of solutions (9) are called the *generating solutions* since these are the first crude approximations one starts with. More elaborate solutions are given by Eqs. (12) and (13) which are termed *guiding centers*.

Notice two points that difference the case considered from the usual averaging techniques. The first point is that in Eq. (8) the small parameter ε is not set zero. And the second difference is in the occurrence of the stochastic average in Eq. (10). Leaving ε in Eq. (8) makes it possible to correctly take into account attenuation effects, as will be shown in applications.

2.4 Generalized expansion about guiding centers

Higher–order corrections to solutions may be obtained by presenting the latter as asymptotic expansions about the guiding centers (12) and (13). To this end, k–order approximations are written as

$$u_k = y(\varepsilon, \xi, t) + \sum_{n=1}^{k} y_n(\varepsilon, \xi, t)\, \varepsilon^n \; ,$$

$$s_k = z(\varepsilon, t) + \sum_{n=1}^{k} z_n(\varepsilon, \xi, t)\, \varepsilon^n \; . \tag{14}$$

Such series are called *generalized asymptotic expansions* [11] since the expansion coefficients depend themselves on the parameter ε. The right–hand sides of Eqs. (3) are to be expanded similarly to Eq. (14) yielding

$$f(\varepsilon, u_k, s_k, \xi, t) \simeq f(\varepsilon, y, z, \xi, t) + \sum_{n=1}^{k} f_n(\varepsilon, \xi, t)\, \varepsilon^n \tag{15}$$

and an equivalent expansion for g. These expansions are to be substituted into Eqs. (3) with equating the like terms with respect to the powers of ε. In the first order, this gives

$$\frac{dy_1}{dt} = f_1(\varepsilon, \xi, t) - \overline{g}(\varepsilon, z) X_1(\varepsilon, \xi, t) \; , \qquad \frac{dz_1}{dt} = g(\varepsilon, y, z, \xi, t) - \overline{g}(\varepsilon, z) \; , \tag{16}$$

where

$$X_1(\varepsilon, \xi, t) \equiv \frac{\partial}{\partial z} X(\varepsilon, z, \xi, t) \; , \qquad z = z(\varepsilon, t) \; .$$

For the approximations of order $n \geq 2$, one gets

$$\frac{dy_n}{dt} = f_n(\varepsilon, \xi, t) \; , \qquad \frac{dz_n}{dt} = g_n(\varepsilon, \xi, t) \; . \tag{17}$$

The functions f_n and g_n depend on $y_1, y_2, \ldots, y_n$ and on $z_1, z_2, \ldots, z_n$ (see for details [6,7]). But it is important that the dependence on y_n and z_n is linear. Therefore all equations (16) and (17) are linear and can be easily integrated. Thus, the approximants (14) are defined. Each k–order approximation can also be improved by invoking the self–similar summation of asymptotic series [12–18].

2.5 Selection of scales for space structures

The solutions of differential or integro–differential equations in partial derivatives are often nonuniform in space exhibiting the formation of different spatial structures. Also, it often happens that a given set of equations possesses several solutions corresponding to different spatial patterns or to different scales of such patterns [3]. When one has a set of solutions describing different possible patterns, the question arises which of these solutions, and respectively patterns, to prefer? This problem of pattern selection is a general and very important problem constantly arising in considering spatial structures. In some cases this problem can be solved as follows.

Assume that the obtained solutions describe spatial structures that can be parametrized by a multiparameter b, so that the k–order approximations $u_k(b, t)$ and $s_k(b, t)$ include the dependence on b whose value is however yet undefined. To define b, and respectively the related pattern, one may proceed in the spirit of the self-similar approximation theory [12–14], by treating b as a control function, or a set of control functions if b is a multiparameter. According to the theory [12–14], control functions are to be defined from fixed–point conditions for an approximation cascade, which is to be constructed for an observed quantity. For the latter, one may take the energy which is a functional $E[u, s]$ of the solutions. In experiments, one usually measures an average energy whose k–order approximation writes

$$E_k(b) \equiv \ll \frac{1}{\tau} \int_0^\tau E[u_k(b,t), s_k(b,t)]\, dt \gg , \tag{18}$$

where τ is a period of fast oscillations. For the sequence of approximations, $\{E_k(b)\}$, it is possible to construct an approximation cascade [12–14] and to show that its fixed point can be given by the condition

$$\frac{\partial}{\partial b} E_k(b) = 0 , \tag{19}$$

from which one gets the control function $b = b_k$ defining the corresponding pattern. According to optimal control theory, control functions are defined so that to minimize a cost functional. In this case, it is natural to take for the latter the average energy (18). Therefore, if the fixed–point equation (19) has several solutions, one may select of them that one which minimizes the cost functional (18),

$$E_k(b_k) = \text{abs} \min_b E_k(b) . \tag{20}$$

Equations (19) and (20) have a simple physical interpretation as the minimum conditions for the average energy (18). However, one should keep in mind that there is no in general such a principle of minimal energy for nonequilibrium systems [3]. Therefore the usage of the ideas from the self–similar approximation theory [12–14] provides a justification for employing conditions (19) and (20) for nonequilibrium processes.

In the following sections a brief survey is given of several physical examples the scale separation approach has been applied to, and the main results are formulated.

3 Superradiance of Nuclear Spins

A system of neutral spins in an external magnetic field , prepared in a strongly nonequilibrium state and coupled with a resonance electric circuit, displays rather nontrivial relaxation behaviour somewhat similar to that of an inverted system of atoms. This is why the optical terminology, such as superradiance, has been used for describing collective relaxation processes in nonequilibrium nuclear magnets [5,6,19,20].

For a system of nuclear spins interacting through dipole forces the evolution equations can be derived [5,6] for the averages

$$u \equiv \frac{1}{N} \sum_{i=1}^{N} < S_i^- > , \qquad s \equiv \frac{1}{N} \sum_{i=1}^{N} < S_i^z > , \tag{21}$$

in which N is the number of spins, angle brackets mean statistical averaging, S_i^- is a lowering spin operator, and S_i^z is the z–component of a spin operator. Following the ideology of the scale separation approach, local fluctuating fields are presented by stochastic variables ξ_0 and ξ. In this way, one comes to the evolution equations for the transverse spin variable

$$\frac{du}{dt} = i(\omega_0 - \xi_0 + i\gamma_2)u - i(\gamma_3 h + \xi)s \tag{22}$$

and the longitudinal average spin

$$\frac{ds}{dt} = \frac{i}{2}(\gamma_3 h + \xi)u^* - \frac{i}{2}(\gamma_3 h + \xi^*)u - \gamma_1(s - \zeta) . \tag{23}$$

It is also convenient to consider the equation

$$\frac{d}{dt}|u|^2 = -2\gamma_2 |u|^2 - i(\gamma_3 h + \xi)su^* + i(\gamma_3 h + \xi^*)su . \tag{24}$$

In equations (22)–(24) dimensionless units are used for the resonator magnetic field h satisfying the Kirchhoff equation

$$\frac{dh}{dt} + 2\gamma_2 h + \omega^2 \int_0^t h(t')\, dt' = -2\alpha_0 \frac{d}{dt}\,(u^* + u) + \gamma_3 f\,. \tag{25}$$

Here ω_0 is the Zeeman frequency of spins in an external uniform magnetic field, ω is the resonator natural frequency, γ_1 and γ_2 are the spin–lattice and spin–spin relaxation parameters, respectively, γ_3 is the resonator ringing width, ζ is a stationary spin polarization, α_0 is the coupling between spins and the resonator, and f is an electromotive force. The random local fields are defined as Gaussian stochastic variables with the stochastic averages

$$\ll \xi_0^2 \gg = \ll |\xi|^2 \gg = \gamma_2^*\,, \tag{26}$$

where γ_2^* is the inhomogeneous dipole broadening.

There are the following small parameters in the system:

$$\frac{\gamma_1}{\omega_0} \ll 1\,, \qquad \frac{\gamma_2}{\omega_0} \ll 1\,, \qquad \frac{\gamma_2^*}{\omega_0} \ll 1\,, \qquad \frac{\gamma_3}{\omega} \ll 1\,,$$

$$\frac{\Delta}{\omega_0} \ll 1\,, \qquad (\Delta \equiv \omega - \omega_0)\,. \tag{27}$$

This makes it admissible to classify the functions u and h as fast, while s and $|u|^2$ as slow, and to apply the method of Section 1. The behaviour of solutions to Eqs. (22)–(25) depends on initial conditions for $u(0)$, and $s(0)$, on the existence of an electromotive driving force $f(t)$, on the pumping related to the parameter ζ, and on the value of the effective coupling parameter

$$g = \pi^2 \eta\, \frac{\rho\, \mu_n^2\, \omega_0}{\hbar\, \gamma_2\, \omega}\,, \tag{28}$$

in which η is a filling factor; ρ, spin density; and μ_n is a nuclear magnetic moment.

The first interesting result is that the electromotive force does not influence much macroscopic samples [5,6] since the corresponding correlation time is proportional to N, that is, the effective, interaction strength of an electromotive force with the spin system is proportional to N^{-1}. This shows, in particular, that the role of the thermal Nyquist noise for starting the relaxation process is negligible. The main cause triggering the motion of spins leading to coherent self-organization is the presence of *nonsecular dipole interactions* [5,6,19]. The latter result gives an answer to the problem, posed by Bloembergen and Pound [21]: What is the origin of self-organized coherent relaxation in spin systems?

All possible regimes of nonlinear spin dynamics have been analysed [5,6,19,20]. When the nonresonant external pumping is absent, that is $\zeta > 0$, there are seven qualitatively different transient relaxation regimes: *free induction, collective induction, free relaxation, collective relaxation, weak superradiance, pure superradiance*, and *triggered superradiance* [6]. In the presence of pumping, realized e.g. by means of dynamical nuclear polarization directing nuclear spins against an external constant magnetic field, one has $\zeta \leq 0$. Then three dynamical regimes can be observed, depending on the value of ζ with respect to the pumping thresholds

$$\zeta_1 = -\frac{1}{g}\,, \qquad \zeta_2 = -\frac{1}{g}\left(1 + \frac{\gamma_1^*}{2\gamma_2}\right)\,, \tag{29}$$

where γ_1^* is an effective pumping rate.

Two stationary points can exist for the slow solutions s and w, where

$$w \equiv |u|^2 - 2\left(\frac{\gamma_2^*}{\omega_0}\right)^2 s^2\,.$$