

# Effect of the bichromatic electromagnetic field on the random molecular aggregate or the quantum spin chain

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The linear response of the molecular aggregate chain with the random distribution of the strength of the neighboring intermolecular interactions to the bichromatic electromagnetic field has been calculated. The presence of the linearly polarized field produces additional resonance absorption at the combined frequencies of the circularly polarized and linear polarized field. The strength of the absorption oscillates with the amplitude and the frequency of the linearly polarized electromagnetic field. Additional resonance absorption is weaker than the main absorption at the frequency of the circularly polarized field. Similar behavior has been predicted for the interaction of the bichromatic magnetic field with the quantum spin-1/2 chain with the random distribution of the exchange constants for nearest neighbor spins.

Keywords: molecular aggregate, bichromatic electromagnetic field, spin chain.

## Introduction

A great interest of researchers has been directed during the last decade to organic nanoscale materials [1]. It is connected with the potential application of those materials in optoelectronics and photonics [2]. Molecules of such systems manifest high fluorescence and absorption efficiencies at low densities and are easy and cheap in applications. In many devices based on organic nanoscale materials (like photovoltaic devices, light-emitting diodes, electrochromic devices, field-effect transistors, etc.) low-dimensional organic systems are used. There it is necessary to prevent interaction between molecules or aggregation because the latter often lead to the fluorescence or absorption quenching due to electronic excitation energy re-distribution [3]. There exist, however, organic fluorophor aggregates which reveal a strong fluorescence [4]. Molecular aggregates are nanoclusters of non-covalently coupled dyes such as cyanines, merocyanides, porphyrins, etc. They were discovered independently by E. Jelley and G. Scheibe [5,6]. The observed features of their optical properties were described by the formation of the dye aggregates. Aggregates are

related to the class of large nanoclusters with the length  $L \geq 100$  nm and the translational symmetry [7]. Due to translational symmetry within the aggregate molecular chains, the electronic excitations of monomers delocalize over chain-forming Frenkel excitons [8]. Usually, one distinguishes so-called *J*-aggregates [9], which demonstrate the long-wavelength (blue) shift of the resonance absorption and fluorescence, with respect to the ones from a single molecule, and *H*-aggregate (called so due to the hypsochromatic shift), which demonstrate the short-wavelength (red) shift. Unlike *J*-aggregates, the absorption line of the *H*-aggregates usually contains a great degree of the vibration effect, and its bandwidth is larger than the width of the absorption in *J*-aggregates (of order of the one for monomers). Some compounds, e.g., the PIC (pseudocyanine) dye can manifest both *J*- and *H*-lines when aggregating [10]. It is often necessary to change (govern) the optical parameters of the dye not changing their chemical structure.

On the other hand, due to the development of spintronics and quantum spin liquids [11] during the recent years, the interest to quantum spin systems has been grown

considerably. Unlike higher-dimensional models, their low-dimensional quantum counterparts permit to obtain exact (non-perturbative) results [12]. The necessity of non-perturbative studies in low-dimensional quantum systems is caused by the enhanced quantum and thermal fluctuations there, due to the features in the density of states [13].

Exactly solvable quantum models are not only important for the pure quantum many-body theory. Nowadays they are extremely useful in a number of areas, which are described by that theory, like real quasi-low-dimensional spin systems, manufactured during the last decade, topological insulators, ultracold atoms in optical traps, strongly correlated electrons interacting via superexchange orbital interactions, sets of qubits (subjects of quantum information and computation), etc. Most of the mentioned issues cannot be explained within the framework of weak couplings (the perturbative theory), or the mean-field-like approach.

In this paper we study theoretically, using the exact quantum solution, the resonance absorption of the bichromatic electromagnetic field by a one-dimensional molecular aggregate with the random distribution of the interaction between neighboring molecules. We show that the absorption can manifest contributions from the combined resonances, which appear due to the coherent interference resonance effects of the bichromatic field. With the help of combined resonances, it is relatively easy to shift the absorption spectrum band to other wavelengths, different from the ones for the standard monochromatic field. On the other hand, we show how the randomness in the distribution of the coupling constants can strongly affect the shape of the resonance absorption spectra. Also, we show that our results can be also applied to the description of the absorption of the ac magnetic field by quantum spin chains with the random distribution of exchange constants there.

### Excitons in molecular aggregates

Excitons are electron-hole pairs, i.e., they carry no electric charge [14]. In molecular aggregates one usually deals with Frenkel excitons, because the charge overlap in neighboring molecules is small [8]: The electron and hole of the Frenkel exciton exist on the same molecule in molecular aggregates. However, excitons are able to propagate through the aggregate as a pair. The main contribution to that propagation comes from interactions between transition dipoles, which form the exciton band. For the ensemble of interacting quantum molecules, one can neglect couplings between ground states and singly and doubly excited states. The interaction between excitons in molecular aggregates is of the Coulomb (electric) nature. It is dominated by the dipole-dipole coupling between molecular transition dipoles. For the description of the ensemble of excitons usually one uses the second quantization description. Within the Hartree–Fock approximation, the excited state of the exciton is formed by the creation of an electron in an unoccupied (excited)

orbital and simultaneous destruction of an electron state of the occupied (ground state) orbital, neglecting intermolecular charge transfer. Such creation and annihilation of the  $n$ th two-level molecule can be described in terms of Pauli creation and destruction operators,  $B_n^\dagger$  and  $B_n$ . The latter satisfy the following commutation relations

$$\begin{aligned} B_n^\dagger B_n + B_n B_n^\dagger &= 1, \quad (B_n)^2 = (B_n^\dagger)^2 = 0, \\ [B_n, B_j] &= 0, \quad [B_n, B_j^\dagger] = 0, \quad n \neq j. \end{aligned} \quad (1)$$

The problem with the quantum mechanical description of such an ensemble of Frenkel excitons is in these commutation relations, which are bosonic for different molecules and fermionic for the same molecule. That produces so-called kinematic interaction between excitons. Using these notations the Hamiltonian of the ensemble of excitons can be written as

$$\begin{aligned} \mathcal{H}_0 &= \sum_n \left[ (\hbar\omega_0 + D_n) B_n^\dagger B_n + \right. \\ &\left. + \sum_j J_{n,j} (B_n^\dagger B_j + B_j^\dagger B_n) \right] + \mathcal{H}_{\text{int}} + E_g, \end{aligned} \quad (2)$$

where  $\hbar\omega_0$  is the energy difference between the ground state and the excited state of the exciton at the site  $n$  of the aggregate,  $J_{n,j}$  is the strength of the interaction between excitons at sites  $n$  and  $j$ , and  $D_n$  is the difference between the ground state in the Coulomb interaction of the first interacting molecule and the excited state of the other interacting molecule, or vice versa (we assume that molecules are situated symmetrically). Finally,  $\mathcal{H}_{\text{int}}$  is the part of the Hamiltonian, which contains higher-order in  $B_n$  and  $B_n^\dagger$  operators ( $\mathcal{H}_{\text{int}}$  is often called the dynamic interaction), and  $E_g$  is the ground-state energy (which expression does not depend on exciton operators). Within the dipole approximation the dynamic interaction part  $\mathcal{H}_{\text{int}}$  can be neglected. However, the kinematic interaction persists, so the quadratic form of Pauli creation and destruction operators still describes the interacting system. The dipole operator can be written, e.g., as

$$\hat{\mu} = \mu \sum_n (B_n^\dagger + B_n). \quad (3)$$

It follows then that in linear optics one can only examine properties of one-exciton states. The other projection of the dipole operator can be written as  $\hat{\mu}_z = \mu_z \sum_n B_n^\dagger B_n + C$ , where  $C$  does not depend on the exciton operators.

### Interaction of excitons with the bichromatic electromagnetic field

Consider now the action of the bichromatic electromagnetic field

$$\mathbf{E}(t) = -E_1(\mathbf{x} \cos(\omega t) + \mathbf{y} \sin(\omega t)) - E_2 \mathbf{z} \cos(\Omega t), \quad (4)$$

i.e., the electromagnetic field has the component with the frequency  $\omega$  and the magnitude  $E_1$ , circular polarized in the  $xy$  plane, and the component with the frequency  $\Omega$  and the magnitude  $E_2$ , linear polarized parallel to the  $z$  axis. Suppose the matrix elements of the dipole operators have the components  $\mu$ , and  $\mu_z$ , perpendicular and parallel to the  $z$  axis, respectively. Then the coupling of the electromagnetic field to the system can be described by the Hamiltonian

$$\mathcal{H}_t = -(\mu E_1 / 2) \sum_n \left( B_n^\dagger e^{-i\omega t} + \text{H.c.} \right) - \mu_z E_2 \cos(\Omega t) \sum_n B_n^\dagger B_n + C(t), \quad (5)$$

where  $C(t)$  is the part, which does not depend on operators. Notice that if  $E_1 = 0$  the molecular system does not respond to the electromagnetic field at all. Let us use the unitary transformation

$$U = \exp \left[ -i\mu_z E_2 \sin(\Omega t) \sum_n B_n^\dagger B_n / \hbar\Omega \right]. \quad (6)$$

The time-dependent part of the considered Hamiltonian can be transformed to

$$\mathcal{H}_t^{(1)} = -(\mu E_1 / 2) \sum_n \left( B_n^\dagger \exp[-i\omega t + i(\mu_z E_2 / \hbar\Omega) \sin(\Omega t)] + \text{H.c.} \right) + C(t). \quad (7)$$

Then we can use the equality

$$\exp[iz \sin(\Omega t)] = \sum_{m=-\infty}^{\infty} J_m(z) \exp(im\Omega t), \quad (8)$$

where  $J_m(z)$  is the Bessel function. Notice that  $J_{-m}(z) = -J_m(z)$ . It follows that

$$\mathcal{H}_t^{(1)} = -(\mu E_1 / 2) \sum_n \sum_{m=-\infty}^{\infty} J_m(\mu_z E_2 / \hbar\Omega) \times \left( B_n^\dagger \exp[-i(\omega - m\Omega)t] + \text{H.c.} \right) + C(t). \quad (9)$$

Then, as usually, we can use the transformation to the rotated frame with the frequency  $\omega - m\Omega$ . The final Hamiltonian has no explicit time dependence. In fact that expression implies that the effect of the bichromatic electromagnetic field is reduced to the effect of the number of circularly polarized fields with the magnitude  $E_1 J_m(\mu_z E_2 / \hbar\Omega)$  and the frequencies  $\omega - m\Omega$ . The linear polarized component of the field produces additional resonances with  $m = \pm 1, \pm 2, \dots$ , with respect to the case of the only circular polarized field, which corresponds to the case  $m = 0$  and  $E_2 = 0$ . The number of possible resonances is determined by the relative

relations between the values  $\omega_0$ ,  $J_n / \hbar$ , and  $\omega$  and  $\Omega$ . The absolute value of the Bessel function is the decaying oscillating function of its argument, hence the absorption and the fluorescence of the aggregate under the action of the bichromatic electromagnetic field must show quasi-periodic modulations as a function of the magnitude and the frequency of the linearly polarized field. The contributions from the resonances with larger values of  $m$  become essential for larger values of  $\mu_z E_2 / \hbar\Omega$ .

### Random molecular aggregate chain

Now we limit ourselves with the one-dimensional quantum molecular aggregates and suppose the interaction  $J_{n,j}$  to be limited to the nearest neighboring molecules. Also we neglect  $D_n$ , i.e., consider equidistant molecules of the molecular aggregates. In such a case we can write the Hamiltonian of the one-dimensional ensemble of excitons as

$$\mathcal{H}_0 = \sum_n \left[ \hbar\omega_0 B_n^\dagger B_n + J_n (B_n^\dagger B_{n+1} + \text{H.c.}) \right]. \quad (10)$$

Then, using the Jordan–Wigner transformation [15]

$$\begin{aligned} B_n^\dagger &= \prod_{m<n} (1 - 2b_m^\dagger b_m) b_n, \\ B_n &= b_n^\dagger \prod_{m<n} (1 - 2b_m^\dagger b_m), \end{aligned} \quad (11)$$

where  $b_n$  ( $b_n^\dagger$ ) are destruction (creation) operators of a fermion in the  $n$ th molecule, we can re-write the Hamiltonian  $\mathcal{H}_0$  as

$$\mathcal{H}_0 = \sum_n \left[ \hbar\omega_0 b_n^\dagger b_n - J_n (b_n^\dagger b_{n+1} + \text{H.c.}) \right]. \quad (12)$$

The advantage of the transition to fermion operators is in the absence of the kinematic interaction, present in the ensemble of Pauli operators. For the homogeneous interaction  $J_n = J$  the Hamiltonian can be easily diagonalized with the help of the Fourier transform. The case  $J > 0$  describes the long-wave shift of the absorption of the aggregate with respect to the monomer absorption (the so-called  $J$ -line), and for  $J < 0$  the shift to the short-range part of the spectrum takes place (the so-called  $H$ -line).

Consider now the absorption of the bichromatic electromagnetic field by the one-dimensional molecular aggregate if  $J_n = x_n |J|$  with  $x_n = \hat{A} \pm 1$ , i.e., the interaction strength between neighboring sites  $n$  and  $n+1$  can be of the  $J$ -type if  $x_n |J| > 0$ , or of the  $H$ -type if  $x_n |J| < 0$ , depending on the configuration of the values  $x_1, \dots, x_L$  ( $L$  is the length of the chain). The gauge transformation

$$b_n \rightarrow c_n = x_1 \dots x_{n-1} b_n, \quad (13)$$

and the analogous one for the creation operator transforms the Hamiltonian of the exciton system of the one-di-

dimensional aggregate with the randomly distributed signs of the interaction strength ( $c_n$  and  $c_n^\dagger$  are also fermionic operators of creation and destruction, because the unitary linear gauge transformation does not change the statistics of quasiparticles) to the expression, in which all operators  $b_n$  ( $b_n^\dagger$ ) are formally changed to  $c_n$  ( $c_n^\dagger$ ) and  $J_n \rightarrow |J|$ . Hence, we have obtained the Hamiltonian of the homogeneous molecular chain, which does not depend on the set of random parameters  $x_1, \dots, x_L$ . Notice, on the other hand, that (dynamical) correlation functions, necessary to describe the interaction of the dipoles of the aggregate chain with the electromagnetic field, do depend on the set of the inhomogeneity parameters  $x_1, \dots, x_L$ .

Consider the distribution of the random parameters [16]

$$p(x_n) = p\delta(x_n + 1) + (1-p)\delta(x_n - 1), \quad (14)$$

where the parameter  $0 \leq p \leq 1$  determines the strength of randomness. We are interested in the average with respect to the randomness expectation values like

$$A_{\text{av}} = \prod_n \int_{-\infty}^{\infty} dx_n p(x_n) A.$$

For the pair correlation functions we have

$$\langle b_n^\dagger(t) b_{n+m} \rangle_{\text{av}} = (1-2p)^{|m|} \langle c_n^\dagger(t) c_{n+m} \rangle. \quad (15)$$

Let us introduce the correlation length  $\xi = -1/\ln|1-2p|$ . Then the pair correlation functions can be written for  $0 \leq p \leq 1/2$  as

$$\langle b_n^\dagger(t) b_{n+m} \rangle_{\text{av}} = \exp(-|m|/\xi) \langle c_n^\dagger(t) c_{n+m} \rangle$$

and for  $1/2 \leq p \leq 1$  as

$$\langle b_n^\dagger(t) b_{n+m} \rangle_{\text{av}} = (-1)^m \exp(-|m|/\xi) \langle c_n^\dagger(t) c_{n+m} \rangle.$$

### Dynamical susceptibility and the absorption of the electromagnetic field

For  $0 \leq p \leq 1/2$  we can obtain for the ground-state imaginary part of the dynamical susceptibility

$$\bar{\chi}''(k, \omega) = \sum_{m=0, \pm 1, \pm 2, \dots} e^{-ikm - |m|/\xi} \int_{-\infty}^{\infty} dt e^{i\omega t} \langle c_n^\dagger c_{n+m} \rangle / 2\hbar. \quad (16)$$

For  $1/2 \leq p \leq 1$  the additional factor  $(-1)^m$  appears. For the absorption of the electromagnetic field, homogeneous along the chain, we can use  $k = 0$  (for both cases). Hence, the averaged with respect to the considered randomness, dynamical susceptibility is reduced to the one of chain of the molecular aggregates with the homogeneous distribution of the interaction strength multiplied by the factor, which takes into account the randomness. Notice that such simple results can be obtained for the ground state only.

For nonzero temperatures, the calculation of the correlation functions is much more complicated. However, for  $\hbar\omega_0 > |J|$  we can apply the results [17], using which it is possible to show that  $\chi''(k, \omega) \sim (\pi/2\hbar)\delta(\hbar\omega - |\hbar\omega_0| - |J|\cos k)$  for the homogeneous chain.

The absorption of the electromagnetic field  $Q$  is defined as  $Q = \langle \partial \mathcal{H}_t / \partial t \rangle_{t, \text{av}}$ , where brackets denote ensemble averaging with the density matrix, and the index  $t, \text{av}$  denotes the time averaging  $A_{t, \text{av}} = \lim_{\tau \rightarrow \infty} \tau^{-1} \int_0^\tau A(t) dt$ . The absorption of the bichromatic electromagnetic field in the linear response regime can be written as

$$Q = \frac{(\mu E_1)^2}{2} \sum_m J_m^2 (\mu_z E_2 / \hbar\Omega) (\omega - m\Omega) \chi(k=0, \omega - m\Omega). \quad (17)$$

For the case  $0 \leq p \leq 1/2$  for the absorption of the bichromatic electromagnetic field by the aggregate chain in the ground state with random distribution of the signs of the interaction strength we obtain for  $\hbar\omega_0 \gg |J|$

$$Q = \sum_m \frac{(\mu E_1)^2 J_m^2 (\mu_z E_2 / \hbar\Omega) (\omega - m\Omega) \sinh(\xi^{-1})}{2\hbar} \times \frac{\Theta(\hbar(\omega - m\Omega - |\omega_0|) - |J|)}{\sqrt{|J|^2 - \hbar^2(\omega - m\Omega - \omega_0)^2}} \times \frac{|J| \Theta(\hbar(|\omega_0| - \omega + m\Omega) + |J|)}{|J| \cosh(\xi^{-1}) - \hbar(\omega - m\Omega - |\omega_0|)}. \quad (18)$$

Here the step Heaviside functions  $\Theta(x)$  limit the absorption in the molecular aggregate by the edges of the existence of excitons in the linear molecular aggregate. For the case  $1/2 \leq p \leq 1$  we obtain

$$Q = \sum_m \frac{(\mu E_1)^2 J_m^2 (\mu_z E_2 / \hbar\Omega) (\omega - m\Omega) \sinh(\xi^{-1})}{2\hbar} \times \frac{\Theta(\hbar(\omega - m\Omega - |\omega_0|) - |J|)}{\sqrt{|J|^2 - \hbar^2(\omega - m\Omega - \omega_0)^2}} \times \frac{|J| (\Theta(\hbar(|\omega_0| - \omega + m\Omega) + |J|))}{|J| \cosh(\xi^{-1}) + \hbar(\omega - m\Omega - |\omega_0|)}. \quad (19)$$

Notice the missprint in Ref. 18 (the multiplier  $\sinh(\xi^{-1})$  was missed in the expressions for  $Q$ ); however in the calculated results for Figs. 1–3 that multiplier was present). The limit  $\xi \rightarrow 0$  corresponds to the totally random system. The opposite limit  $\xi \rightarrow \infty$  describes the cases with  $p = 0$  and  $p = 1$  (totally homogeneous system). Both limits can be obtained from the expression for the inhomogeneous (for all  $k$ ) part of the dynamical susceptibility (the order of limitation  $k \rightarrow 0$  and  $\xi \rightarrow 0$  or  $\xi \rightarrow \infty$  matters).



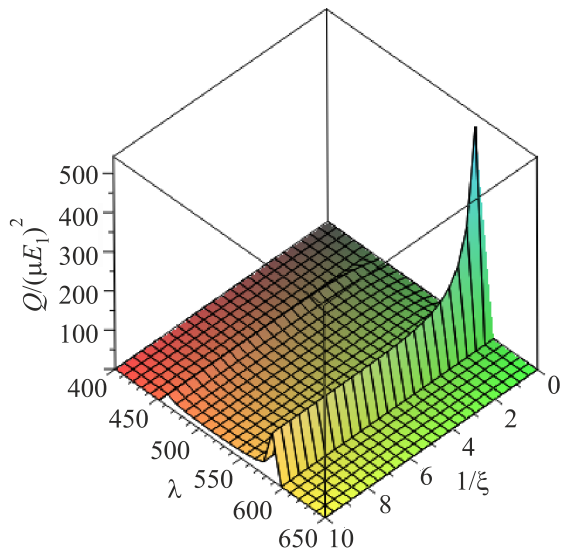


Fig. 1. (Color online) The main contribution  $m = 0$  of the absorption of the bichromatic electromagnetic field by the linear molecular aggregate as a function of the degree of randomness for small number of  $H$ -bonds.

#### Analysis of the obtained results

Consider now the behavior of the ground-state absorption of the bichromatic electromagnetic field by the molecular aggregate chain with the studied randomness. Figure 1 presents the wavelength dependence  $\lambda = 2\pi c/\omega$  ( $c$  is the velocity of light) for the main resonance  $m = 0$  for the  $J$ -aggregate chain with the small degree of  $H$ -aggregate bonds  $0 \leq \xi^{-1} \leq 10$ . The parameters  $\hbar\omega_0 = 0.39$  and  $J = 0.05$  (relative units) were used to obtain the similarity of the calculated results with the ones observed for PIC [19]

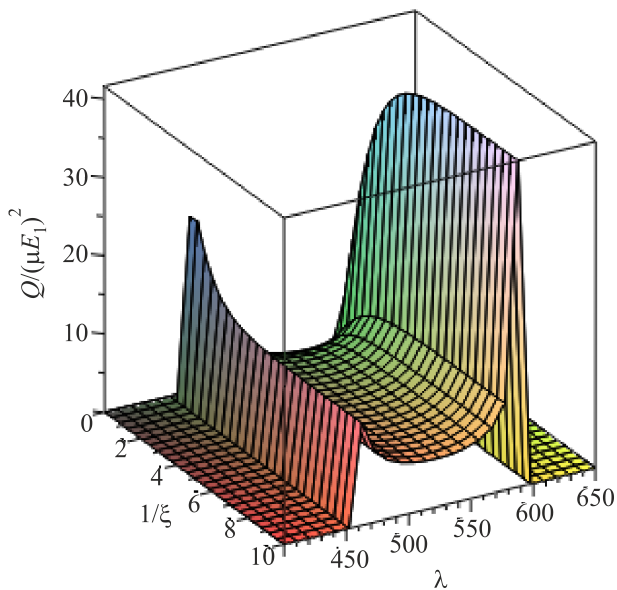


Fig. 3. (Color online) The main contribution  $m = 0$  of the absorption of the bichromatic electromagnetic field by the linear molecular aggregate as a function of the degree of randomness for small number of  $J$ -bonds.

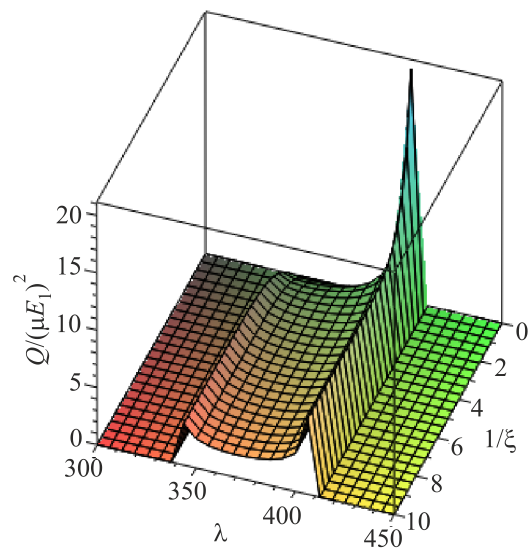


Fig. 2. (Color online) The contribution of the first combined resonance for the same situation as in Fig. 1.

for  $\xi^{-1} = 0$ , to get the long-wave shift, with respect to absorption of the single molecule of PIC. The results, naturally, coincide with the ones of [18]: For the almost homogeneous  $J$ -aggregate chain we see the blue shift, while for smaller values of  $\xi$  the contribution from  $H$ -bonds becomes sharper. On the other hand, in Fig. 2 the results of calculation for the first combined resonance with  $m = 1$  with  $\mu_z E_2 = 0.1$  and  $\hbar\Omega = 0.15$  for the same molecular aggregate chain as in Fig. 1. We see that the absorption at the combined resonance is smaller than the one of the main resonance and takes place for smaller values of the wavelength.

In Figs. 3 and 4 we see the main and the first combined resonances, respectively, for the absorption of the

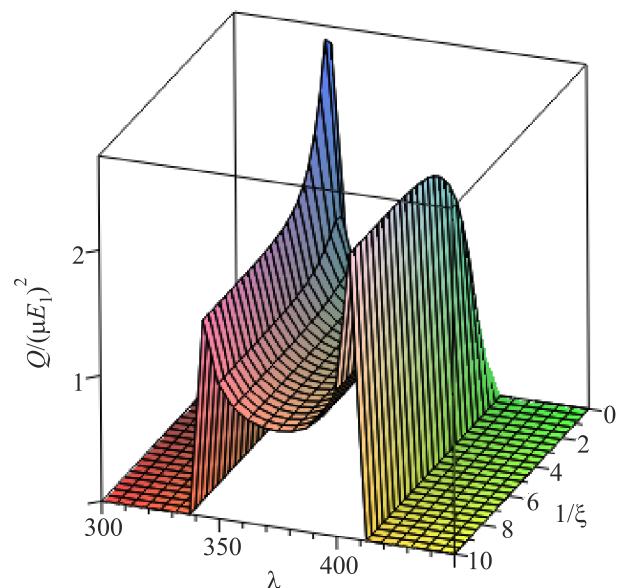


Fig. 4. (Color online) The contribution of the first combined resonance for the same situation as in Fig. 3.

bichromatic electromagnetic field by the  $H$ -aggregate chain with the small degree of  $J$  bonds. For the almost homogeneous case we see the red shift, comparing to the single molecule case. Also we see that the value of the absorption in that case is smaller than for the  $J$ -aggregate (it agrees with the data of experiments on molecular aggregates). The addition of  $J$ -bonds yields the re-distribution of the spectral weight to the blue part of the absorption spectrum. For the combined resonance the absorption is shifted to the short-wavelength region, and it is smaller than the main resonance. Comparing to the previous case, we see that the relative weights of the  $H$  and  $j$  wings for small finite values of the degree of randomness,  $\xi^{-1}$ , are different for the main and the first combined resonance absorption.

We notice also that for the combined resonance the larger values of absorption at the middle of the zone with respect to the edges are characteristic, unlike the case of the main resonance.

### Spin-1/2 chain

It is interesting to note that the same results are obtained for the spin-1/2 chain [12] with the Hamiltonian

$$\mathcal{H}_0 = \sum_n \left[ -g_z \mu_B H S_n^z - J_n (S_n^x S_{n+1}^x + S_n^y S_{n+1}^y) \right], \quad (20)$$

where  $S_n^{x,y,z}$  are the operators of the projections of the spin at the site  $n$ ,  $\mu_B$  is the Bohr magneton,  $g_z$  is the effective  $g$ -factor along  $z$  direction,  $H$  is the external dc magnetic field, and  $J_n$  are the exchange integrals between neighboring spins of the chain, interacting with the ac bichromatic magnetic field

$$\begin{aligned} \mathcal{H}_t = & -(g \mu_B h_1 / 2) \sum_n (S_n^- e^{-i\omega t} + \text{H.c.}) - \\ & - g_z \mu_B h_2 \cos(\Omega t) \sum_n S_n^z + C(t), \end{aligned} \quad (21)$$

where  $g$  is the  $g$ -factor along  $x$  and  $y$  directions,  $h_{1,2}$  are the magnitudes, and  $\omega$  and  $\Omega$  are the frequencies of the circularly polarized and linearly polarized ac magnetic field, respectively. Using the Jordan–Wigner transformation for spin-1/2 operators, we can show that up to re-notation  $\hbar\omega_0 \rightarrow g_z \mu_B H$ ,  $\mu E_1 \rightarrow g \mu_B h_1$  and  $\mu_z E_2 \rightarrow g_z \mu_B h_2$  the obtained above results can be applied for description of the absorption of the bichromatic ac magnetic field by the quantum spin-1/2 chain with the randomly distributed signs of the exchange coupling between neighboring spins. Notice, that for the spin chain the resonance absorption takes place for smaller frequencies (rf and far infrared region, comparing to the optical diapason for molecular aggregates).

### Summary

Summarizing, we have calculated the linear response of the linear molecular aggregate with the random distribution

of the strength of the neighboring intermolecular interactions to the bichromatic electromagnetic field. It has been shown that the presence of the linearly polarized field produced additional resonance absorption at the combined frequencies of the circularly polarized and linear polarized fields. The strength of the absorption oscillates with the amplitude and the frequency of the linearly polarized electromagnetic field. Additional resonance absorption is weaker than the main absorption at the frequency of the circularly polarized field. For combined resonances the relative weights of the short-wavelength and long-wavelength wings of the absorption are different from the ones of the main resonance (only circularly polarized field). We have also shown that even small concentration of the impurities with the different sign of the interaction between molecules of the aggregate chain can strongly renormalize the absorption spectrum. On the other hand, for the strong degree of randomness small changes of the degree of disorder does not change drastically the character of the absorption. We have shown that similar behavior is predicted for the interaction of the bichromatic ac magnetic field with the spin-1/2 chain with the random distribution of the exchange constants for nearest neighbor spins.

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Вплив біхроматичного електромагнітного поля  
на неупорядкований молекулярний агрегат  
або квантовий спіновий ланцюжок

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Розраховано лінійний відгук ланцюжка молекулярних агрегатів із випадковим розподілом сили сусідніх міжмолекулярних взаємодій на біхроматичне електромагнітне поле. Наявність лінійно поляризованого поля викликає додаткове резонансне поглинання на комбінованих частотах з круговою поляризацією по лінійній поляризації поля. Сила поглинання коливається з амплітудою та частотою лінійно поляризованого електромагнітного поля. Додаткове резонансне поглинання слабкіше основного поглинання з круговою поляризацією. Аналогічну поведінку передбачено для взаємодії біхроматичного магнітного поля з квантовим ланцюжком спінів  $1/2$  з випадковим розподілом обмінних констант для найближчих сусідніх спінів.

Ключові слова: молекулярний агрегат, біхроматичне електромагнітне поле, спіновий ланцюжок.