Conservation of the phase correlation of optical fields during scattering

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An experimental investigation was made of the influence of scattering of wide-band optical fields on the conservation of their phase correlation. The correlation was conserved when correlated fields were scattered by bulk or surface scatterers. The experimental results were analyzed theoretically.

INTRODUCTION

The effectiveness of one-photon interaction of wide-band fields with resonant systems depends weakly on the statistics of the optical radiation, but is governed mainly by the spectral density of the radiation power in an absorption band (profile) of a resonant system. As pointed out in a general form in Ref. 1, in the case of many-photon processes this description becomes generally invalid and statistical properties of the incident radiation may determine largely the nature and effectiveness of the nonlinear optical phenomena induced by this radiation. In particular, it was shown in Ref. 2 that wide-band radiation with a specific pair correlation of harmonics can interact with a resonant system in the same way as a quasimonochromatic field with a fixed frequency ω_0. When the frequency $2\omega_0$ coincides with the natural frequency of a resonant system, the effectiveness of two-photon excitation of the system is maximal and is governed not by the spectral power density within the resonance profile, but by the total intensity of the field. Wide-band radiation can be regarded as a set of pairwise coupled harmonics with frequencies and phases satisfying the conditions $\omega+\omega'=2\omega_0$ and $\varphi_\omega+\varphi_{\omega'}=\text{const}$. The first relationship represents the usual condition for a resonance of a two-photon transition, whereas the second requires the sums of the phases of all the pairwise coupled harmonics to be the same but the phases themselves can be random quantities. This condition ensures maximum "monochromatization" of two-photon excitation of a resonant system by a wide-band field.

Other variants of the organization of the field with pairwise correlation of the components, leading to monochromatization of the interaction, are also possible. For example, it is suggested in Ref. 3 that coherent harmonics, shifted by a frequency ω_0, can be generated from the original wide-band radiation. The original and new harmonics with frequencies related by $\omega-\omega'=\omega_0$ obey the phase relationships $\varphi_\omega-\varphi_{\omega'}=\text{const}$. The simultaneous action of the original and frequency-shifted wide-band signals gives rise to a quasimonochromatic effect in four-photon parametric oscillation in a three-level system.

The problem of stability of intraspectral correlation in the course of propagation of radiation is important in the case of correlated optical fields. The most obvious decorrelating factors of practical importance are the frequency dispersion of the refractive index and the inhomogeneity of the medium resulting in optical radiation scattering.

The dispersion-induced decorrelation can easily be estimated on the basis of the following ideas. Total decorrelation corresponds to a uniform distribution of the sum of the phases $\varphi_{\omega_0+\nu}+\varphi_{\omega_0-\nu}$ of different pairs of coupled harmonics in the interval 2π ($\omega_0+\nu=\omega, \omega_0-\nu=\omega'$). After traveling a distance l in the investigated medium, the net phase of a pair is shifted by

$$\delta_\nu = k_0 l \nu^2 \left(\frac{2}{\omega_0}\frac{\partial n}{\partial \omega_0}+\frac{\partial^2 n}{\partial \omega_0^2}\right),$$

where $k_0=\omega_0/c$ is the wave number of the central harmonic and n is the refractive index. It is clear from this expression that the greater the separation of the coupled harmonics from the center $\nu=0$, the smaller the value of l at which the phase shift amounts to 2π relative to the central harmonic. For the same values of l the central part of the spectrum which satisfies the condition $\delta_\nu \ll 2\pi$ remains correlated.

It is much more difficult to describe the loss of correlation in the course of scattering of radiation by inhomogeneities in a medium and such a description depends strongly on the actual properties of the scattering medium. We can only say definitely that the intraspectral correlation is lost as a result of scattering if different pairs of coupled harmonics undergo very different phase distortions. Consequently, it is natural to expect that in the case of strong scattering, when the front of the incident wave is deformed regularly and strongly and spatial coherence is lost and there is also loss of intraspectral correlation of the field. We shall show below that in the case of such scattering the initial correlation may be largely conserved. This experimentally observed effect is attributed to the specific properties of the stability of a fourth-order correlation function of the field which ensures the monochromatic effect in two-photon excitation of resonant systems for the scatterers used in our experiments.

EXPERIMENTS

The main purpose of our experiments was to investigate changes in intraspectral correlation of harmonics of stochastic radiation experiencing surface and bulk scattering. We used a method developed earlier for the generation of fields with a frequency-phase correlation of harmonics in the optical range, two-photon excitation of a cesium atomic vapor, and detection of the luminescence emitted by excited atoms.[4] Figure 1 shows schematically the apparatus. Multimode radiation emitted by a dye (polymethine 1605) laser (components 2–8 in Fig. 1), used together with an amplifier 9 and pumped by single pulses from a ruby laser 1, was the original wide-band radiation with an average wavelength of 930 nm and a discretely variable width of the spectrum. This width was varied by introducing various selectors (Fabry-Perot interferometers and a diffraction grating) into the cavity, and by combining these we were able to obtain spectral widths 0.5, 3, 10, 23, and 90 cm^{-1}, i.e., we could vary this

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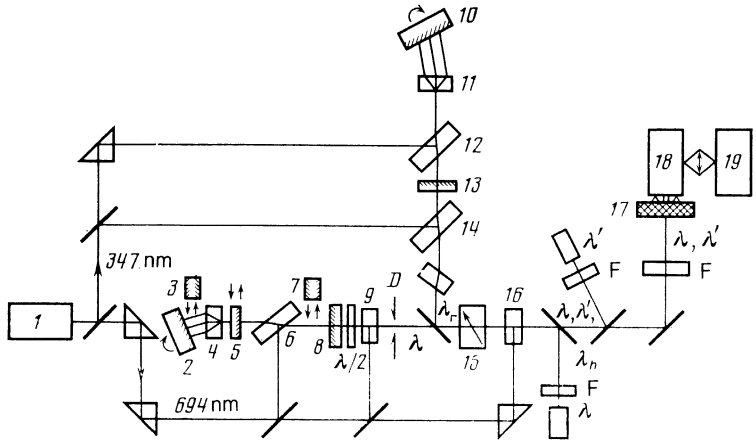

FIG. 1. Diagram of the apparatus: 1) source emitting at $\lambda = 694$ and 347 nm; 2), 10) diffraction gratings; 3), 7) Fabry-Perot interferometer; 4), 11) prism beam expanders; 5) mirror with $R = 100\%$; 6), 9), 12), 14), 16) cells with dyes; 8), 13) exit mirrors; 15) nonlinear $LiIO_3$ crystal; 17) scatterers; 18) cell with cesium vapor; 19) monochromator and photomultiplier.

width by a factor of 180. The constancy of the geometry of propagation of the rays was monitored specially so as to avoid the influence of the geometric factors on the degree of correlation of the fields. Narrow-band heterodyne radiation was generated also by a coumarin dye laser (10–13) operating in conjuction with an amplifier 14. This dye laser was pumped by the second harmonic of the ruby laser 1 ($\lambda = 347$ nm). The radiation from this heterodyne laser had a wavelength $\lambda_h = 441.7$ nm and the linewidth was 0.5 cm^{-1}. The wide-band radiations of the first dye laser and of the second (heterodyne) laser were directed to an $LiIO_3$ crystal 15 where they were mixed nonlinearly giving rise to wide-band radiation for which the frequency ω' and phase $\varphi_{\omega'}$ of each harmonic was related to the frequency and phase of the original noise field by $\omega + \omega' = 2\omega_0$ and $\varphi_\omega + \varphi_{\omega'} = \text{const}$. This radiation, together with that which passed through the crystal, formed a stochastic field with a frequency-phase correlation of the harmonics.

Both radiations traveling collinearly (in accordance with the mixing scheme of the wide-band and heterodyne radiations) were directed to a cell 18 containing cesium vapor in a buffer gas atmosphere (He at a pressure of 50 Torr), where the Cs atoms were subjected to two-photon excitation (Fig. 2) in accordance with the transition $6S_{1/2} \rightarrow 6P_{3/2} \rightarrow 6D_{5/2}$ (the offset from the intermediate level was about 80 Å). Efficient two-photon excitation of the atoms required that the intensities of the two fields I and I' should be comparable. Since after passing through the nonlinear crystal in the region of spatially overlapping rays the ratio I/I' was of the order of 50, we placed a dye amplifier 16 (with the gain maximum at 840 nm and pumped by the ruby laser) between the crystal and the cesium cell. Beyond this amplifier the ratio I/I' had an acceptable value of 10:1. The excitation of the Cs atoms was detected by observing the luminescence generated in accordance with the level scheme shown in Fig. 2. The cesium atoms excited to the $6D_{5/2}$ state collided with the He atoms and were thus transferred to the $7P_{3/2,1/2}$ state from which they relaxed to the ground state emitting at $\lambda = 4555$ and 4593 Å. Therefore, it was possible to achieve an effective optical decoupling of the excitation (infrared) and detection (visible) channels. The luminescence from the middle part of the cell was directed via a monochromator 19 to a photomultiplier. Filters F in front of the cell removed the heterodyne radiation and ensured linear operation of the photomultiplier. The criterion of the frequency-phase correlation of the optical field harmonics was theoretically predicted and experimentally confirmed[2,4] by the observation that the efficiency of the two-photon excitation of the cesium atoms was independent of the width of the spectrum of the original noise radiation (Δ_n) when its intensity was constant.

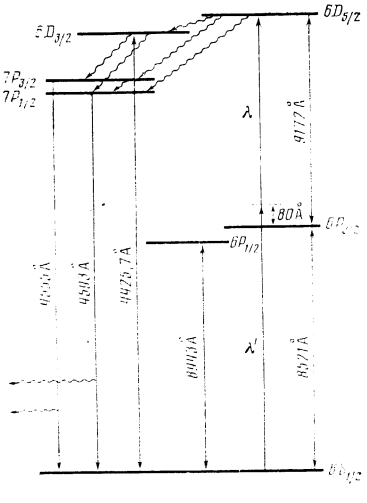

FIG. 2. Energy level scheme of the cesium atom.

Directly in front of the entry window of the cesium cell we placed various scatterers in the path of the radiation characterized by a frequency-phase correlation of the harmonics. It was assumed that the setting of a scatterer resulted in the loss of the phase correlation field because of the destruction of the initial wave front which could occur in different ways for the fields I and I'. We therefore expected that in the case of the noncoherent part of the radiation transmitted by a scatterer the efficiency of the excitation of the Cs atoms should depend on the width of the spectrum of the initial radiation.

We used bulk and surface scatterers. The bulk scattering media were MS-19 light-scattering glass plates[5] (scatterers Nos. 1 and 2). This glass contained inclusions with a characteristic size of ~ 1 μm representing inhomogeneities of the refractive index and distributed homogeneously throughout the bulk of the glass. By altering the thickness of

the scattering glass plate (scatterers Nos. 1 and 2 differed specifically in respect of the thickness), we could alter the parameters of the transmitted light in a controlled manner. The surface scatterers were mat quartz plates with 40 μm characteristic dimensions of uniform surface irregularities (scatterer No. 3), and with very nonuniform irregularities of dimensions exceeding 40 μm (scatterer No. 4). In the case of all these scatterers the noncoherent part of the scattered radiation exceeded the coherent part in absolute magnitude and power density. Special measurements showed that the ratio of the coherent part of the scattered radiation to the intensity of the incident radiation was 0.016, 0.003, 0.02, and 0.0 for the scatterers Nos. 1, 2, 3, and 4, respectively. Therefore, after passing through a scatterer the radiation was almost completely diffuse and the signal representing the two-photon excitation of the Cs atoms, due to the coherent part of the wave, should be attenuated by a factor exceeding 10^3 in the scatterer (proportionally to the square of the reduction in the intensity of the coherent radiation).

Insertion of a scatterer reduced the intensity of the two-photon-excited luminescence of Cs because of a reduction in the power density of the radiation since the cross section of the light flux and the absorption increased in the scatterer. Special calibration measurements indicated that this reduction was a factor of 10–100, depending on the scatterer. We allowed for this effect by normalization of all the experimental results to the signal obtained for narrow-band radiation characterized by $\Delta_n = 0.5$ cm^{-1}.

The correlation properties of the scattered signals were investigated by making control measurements of the excitation of the luminescence in accordance with the level scheme described above using uncorrelated fields (i.e., with the heterodyne laser switched off). It was found that broadening of the spectrum of the original radiation from 3 to 90 cm^{-1} while conserving its intensity reduced the luminescence by a factor of about 40, in satisfactory agreement with the theoretical predictions.

The experimental values of the signals S representing the two-photon-excited luminescence of the cesium atoms in scattered correlated fields with different spectral widths Δ_n were averaged for several series of measurements (Table I). The influence of fluctuations of the radiation intensity on the results of the measurements was avoided by referring the measured value of S to the product II' ($Z = S/II'$). Moreover, the results for each scatterer were normalized for the reasons mentioned above to the signal with $\Delta_n = 0.5$ cm^{-1} [$Z(0.5)$].

The absolute values of S were several orders of magnitude higher than the luminescence intensity which could be excited by the residual coherent part of the light flux beyond the scatterer. We therefore concluded that the values of $Z(\Delta_n)$ listed in Table I were associated with the noncoherent part of the scattered radiation. Hence, the efficiency of the two-photon excitation of the cesium atoms was practically independent of the width of the noise field spectrum. In turn, this indicated that the intraspectral correlation was retained by the field after the strong scattering of the light flux by the bulk and surface scatterers which removed almost completely the coherent part of the wave (and thus destroyed almost completely the plane wavefront of the original radiation).

The results listed in Table I showed also that the efficiency of the two-photon excitation by diffuse radiation decreased on increase in Δ_n. This reduction in the efficiency exceeded the experimental error (50%). We were unable to explain this observation convincingly. It could be due to a difference between the divergences of the original and correlated (formed in the process of mixing) rays with different values of Δ_n, which was enhanced by the scattering.

DISCUSSION OF THE EXPERIMENTAL RESULTS

Interpretation of these unexpected experimental results (because the relative changes of the harmonics in the investigated scattering were large) required identification of the correlation functions of the field responsible for the monochromatization effect. This should make it possible to determine in general the feasibility and conditions for conservation of the intraspectral correlation of the light flux in the case of strong diffuse scattering, accompanied by deformation of the incident wavefront and destruction of the spatial coherence.

The loss of coherence requires large relative changes of the net phases of different pairs of coupled harmonics. The fact that the efficiency of the two-photon excitation of the investigated resonant system is independent of the spectral width, for a given intensity of the radiation, is clear evidence of conservation of the intraspectral correlation in the course of strong scattering.

In a theoretical analysis of the problems of the scattering of correlated fields we shall assume that a two-level system excited weakly by two-photon absorption is used as an optical radiation detector. The detected excitation of this system (i.e., the population ρ of the upper level) is then described by

$$\dot{\rho} = -\gamma_1\rho + V^2(t)\int_{-\infty}^{t} V^{*2}(t')\exp[i(t-t')(\omega - 2\omega_0 + i\gamma_2)] \times dt' + \text{c.c.}, \quad (1)$$

where γ_1 and γ_2 are respectively the relaxation rates of the

TABLE I.

Δ_n, cm^{-1}	$Z(\Delta_n)/Z(0.5)$				
	no scatterer	scatterer			
		№ 1	№ 2	№ 3	№ 4
0.5	1	1	1	1	1
3	1	1	1	1	1
10	1	1	0.8	1	0.9
23	0.9	0.6	0.8	0.4	0.6
90	0.9	0.4	0.8	0.4	0.4

population and the coherence of the transition characterized by a frequency ω_{21}. The quantity V and its complex conjugate V^* are related to the intensity of a field

$$F(\mathbf{r}, t)\exp(-i\omega_0 t) + \text{c.c.},$$

which is regarded as scalar, by the expression $V = AF$, where A is a coefficient proportional to the dipole matrix elements of the transitions.

The amplitude of the scattered field $F(\mathbf{r}, t)$ is related linearly to the amplitude of the incident field $F^0(\mathbf{r}, t)$, which is a set of almost plane waves with a dense spectrum in the vicinity of ω_0, i.e.,

$$F^0(\mathbf{r}, t) = \sum_\nu f_\nu^0(\mathbf{r})\exp(ik_\nu z - i\nu t + i\varphi_\nu), \tag{2}$$

where $k_\nu c = \omega_0 + \nu$ and $\nu \ll \omega_0$, and $f^0(\mathbf{r})$ is a smooth function over a scale of k_ν^{-1}. The initial phases φ_ν are regarded as random quantities distributed uniformly in the interval 2π and related by the condition $\varphi_\nu = -\varphi_{-\nu}$.

Scattering by immobile inhomogeneities causes a transformation $f_\nu^0(\mathbf{r}) \to f_\nu(\mathbf{r})$ governing the form of the amplitude of the scattered field

$$F(\mathbf{r}, t) = \sum_\nu f_\nu(\mathbf{r})\exp(-i\nu t + i\varphi_\nu). \tag{3}$$

This field and its action are stochastic. The experiments record the steady-state average value $\bar{\rho}$, which is given by the expression

$$\gamma_1\rho = \int_0^\infty \langle V^2(\tau)V^{*2}(0)\rangle \exp[-\tau(\gamma_2 - i\Delta)]\,d\tau + \text{c.c.}, \tag{4}$$

where $\Delta = \omega_{21} - 2\omega_0$ and $\langle\ldots\rangle$ represents averaging over the initial phases. Bearing in mind that

$$\langle\exp(i\varphi_\nu + i\varphi_\mu)\rangle = \delta_{\nu,-\mu}$$

and

$$\langle\exp(i\varphi_\nu - i\varphi_\mu)\rangle = \delta_{\nu,\mu},$$

where $\delta_{\alpha,\beta}$ is the Kronecker delta, we find that the correlation function of Eq. (2) is described by

$$A^{-4}\langle V^2(\tau)V^{*2}(0)\rangle = \left|\sum_\nu f_\nu f_\nu^* e^{-i\nu\tau}\right|^2 + \left|\sum_\nu f_\nu f_{-\nu}\right|^2, \tag{5}$$

where the first term on the right makes a contribution to Eq. (2), which is inversely proportional to the width of the spectrum and is regarded as unimportant, whereas the second governs the monochromatic effect. Finally, in the case of the experimentally recorded population under the influence of the scattered field, we obtain an expression

$$\bar{\rho} = 2A^4\gamma_2\left|\sum_\nu f_\nu f_{-\nu}\right|^2 \Big/ \gamma_1(\gamma_2^2 + \Delta^2). \tag{6}$$

It is clear from Eq. (6) that a reduction in the excitation efficiency because of the scattering is associated with dispersion of the phases of the harmonics $f_\nu f_{-\nu}$ amounting to

$$Q = \left|\sum_\nu f_\nu f_{-\nu}\right|^2. \tag{7}$$

In the propagation of radiation without scattering and dispersion there is no decorrelation and the value of Q is not affected:

$$Q(z) = \left|\sum f_\nu^0\exp(ik_\nu z) f_{-\nu}^0\exp(ik_{-\nu}z)\right|^2$$

$$= \left|\sum f_\nu^0 f_{-\nu}^0\right|^2 = Q(0).$$

We shall consider the change in Q as a result of the scattering realized experimentally for two types of scatterers: a) with inclusions of $\sim 1\,\mu$m size in the bulk; b) with a rough surface characterized by an inhomogeneity scale $\sim 40\,\mu$m.

In the former case each inclusion scatters the incident field attenuated by the previous scattering. The multiple scattering effects can be regarded as unimportant, since the angular dispersion of the scattering by individual dielectric particles of size of the order of the wavelength of light is strong (as found experimentally) and exceeds the ratio of the size of the illuminated spot to the thickness of the scatterer. The scattered field of each harmonic at the position of a detector R is of the form

$$f_\nu(\mathbf{R}) = R^{-1}\sum_j \exp(ik|\mathbf{R} - \mathbf{r}_j|)\,b(\mathbf{k}_\nu, \mathbf{r}_j)\exp(ik_\nu z_j), \tag{8}$$

where $\mathbf{r}_j$ is the position of the jth inclusion and $b(\mathbf{k}_\nu, \mathbf{r}_j)$ is the amplitude of the scattering by this inclusion proportional to the amplitude of the corresponding harmonic at a point $\mathbf{r}_j$ of the incident field and dependent on the size and shape of the inclusion and on the wavelength and direction of the scattered radiation, i.e., on $\mathbf{k}_\nu$. Substituting Eq. (8) into Eq. (7) describing Q, we obtain

$$Q = R^{-4}\sum_{\nu,\mu,j,l,m,n} \exp(ik_\nu L_j + ik_{-\nu}L_l - ik_\mu L_m - ik_{-\mu}L_n)$$

$$\times b_{\nu,j}b_{-\nu,l}b_{\mu,m}^* b_{-\mu,n}^*, \tag{9}$$

where for the sake of brevity we introduced $L_s = |\mathbf{R} - \mathbf{r}_s| - z_s$ and $b_{\omega,s} = b(\mathbf{k}_\omega, \mathbf{r}_s)$.

If we assume that the positions of the inclusions are statistically independent and that the interference due to the different inclusions is unimportant in the scattered fields, i.e., if $n = m = l = j$, we obtain the following expression:

$$Q \sim R^{-4}\sum_j \left|\sum_\nu b_{\nu,j}b_{-\nu,j}\right|^2. \tag{10}$$

There are no concrete data on the values of $b_{\nu,j}$. However, we can say that in the case of these inclusions these are smooth functions of the scattering angle varying on a scale of ~ 1. Their dependences on the frequency ν are determined by the frequency dispersion of the inclusion material and by the diffraction effects which far from resonances (as in the present case) vary on a scale greater than ω_0 and are almost constant for the harmonic spectrum $\{\nu\}$. If we ignore these weak dependences on ν, we are left only with the dependence $b_{\nu,j} \propto f_\nu^0(\mathbf{r}_j)$, which gives rise to the same dispersion of the phases both for $b_{\nu,j}b_{-\nu,j}$ and $f_\nu^0 f_{-\nu}^0$. It follows that in this

case there is no decorrelation in the case of strong scattering due to a large angular dispersion $b(\mathbf{k}_\nu, \mathbf{r}_j)$. It should be noted that the dependence of $b_{\nu,j}$ on $\mathbf{r}_j$ is governed by the distribution of the characteristics of the inclusions and by the attenuation of the incident radiation as a result of the scattering process. Under the selected conditions these factors can be regarded as independent of ν and, therefore, unimportant from the point of view of decorrelation.

We shall now consider the case of scattering by a rough surface. As in the first case, we observe a wide angular spectrum indicating modulation of the surface relief. Scattering by large inhomogeneities can be described in two stages. First, using the Snell-Fresnel expressions, we can describe it as refraction by piecewise-flat regions and then, having determined the wave field on the surface, we can apply the methods of the theory of diffraction to describe the field far from the illuminated spot. Naturally, this approach ignores the processes of second scattering and the diffraction effects which appear on reflection or refraction by nonplanar surfaces.

The main feature of the first stage of the description is the independence of the nature of the refraction process of the radiation frequency if we ignore the frequency dispersion of the refractive index. This means that collinear wave vectors of the incident harmonics are transformed into new collinear vectors. The amplitudes of the different harmonics then change in a similar manner. These properties can be described for the transmitted field written down in the form $U_\nu(\mathbf{x})\exp(i\varphi_\nu(\mathbf{x}))$, where $\mathbf{x}$ are the coordinates of the surface, by the expansion

$$\varphi_\nu(\mathbf{x}) \approx \varphi_\nu(\mathbf{x}_0) + k_\nu \mathbf{n}(\mathbf{x}_0)\,(\mathbf{x}-\mathbf{x}_0) \tag{11}$$

in the vicinity of an arbitrary value of $\mathbf{x}_0$ with $\mathbf{n}(\mathbf{x}_0)$, which is the same unit vector representing the direction of the optical field of the light flux in all cases, and in the form

$$U_\nu(\mathbf{x}) = D(\mathbf{x}) f_\nu^0(\mathbf{x}), \tag{12}$$

where $D(\mathbf{x})$ is independent of ν. Next, following the Fresnel-Kirchhoff theory, we find that the field of a harmonic at a detector located at $\mathbf{R}$ far from the illuminated spot is given by the expression

$$f_\nu(\mathbf{R}) \propto R^{-1}\exp(ik_\nu R)\int d\sigma \cos(\measuredangle\, \mathbf{n}(\mathbf{x})\,\mathbf{R})\, U_\nu(\mathbf{x}) \times \exp(i\varphi_\nu(\mathbf{x}) - i\mathbf{k}_\nu\mathbf{x} + ik_\nu\mathbf{x}^2/2R), \tag{13}$$

where $\mathbf{k}_\nu = k_\nu\mathbf{R}/R$. The main contribution to the integral comes from the regions of maximum compensation of the terms $i\mathbf{k}_\nu\cdot\mathbf{x}$ and $\varphi_\nu(\mathbf{x})$, giving rise to rapid oscillations of the integrand. It follows from the property described by Eq. (11), that such compensation can be the same for all values of ν directly in the vicinity of some points $\mathbf{r}_j$. The compensation defect

$$\varphi_\nu(\mathbf{x}) - \varphi_\nu(\mathbf{x}_j) - \mathbf{k}_\nu\mathbf{x}_j \approx k_\nu a^{-1}(\mathbf{x}-\mathbf{x}_j)^2,$$

where a is the size of the roughness inhomogeneity, determines the area $k^{-1}a$ of the contribution in the vicinity of $\mathbf{x}_j$ within which all the other integrands vary only slightly. The vicinities of all the points $\mathbf{x}_j$ give rise to the resultant harmonic field

$$f_\nu(R) \propto R^{-1}\exp(ik_\nu R)\sum_j \exp(i\varphi_\nu(\mathbf{x}_j))\, g_\nu(\mathbf{x}_j), \tag{14}$$

where $g(\mathbf{x}_j)$ are the amplitudes of the local sources or "projectors" emitting light in the direction of the detector.

The positions $\mathbf{x}_j$ can naturally be regarded as random so that the quantity Q can be determined as the average

$$\overline{Q} = \left\langle \left| \sum_\nu f_\nu f_{-\nu} \right|^2 \right\rangle = \sum_{\nu,\mu} \langle f_\nu f_{-\nu} f_\mu^* f_{-\mu}^* \rangle \tag{15}$$

over the positions of the points x_j. Substituting Eqs. (11) and (15) and assuming x_j to be statistically independent, we obtain

$$Q \propto R^{-4} N \sum_{\nu,\mu} \langle \exp\{i\varphi_\nu(\mathbf{x}) + i\varphi_{-\nu}(\mathbf{x}) - i\varphi_\mu(\mathbf{x}) - i\varphi_{-\mu}(\mathbf{x})\} \times g_\nu(\mathbf{x})\, g_{-\nu}(\mathbf{x})\, g_\mu^*(\mathbf{x})\, g_{-\mu}^*(\mathbf{x}) \rangle, \tag{16}$$

where N is the number of local projectors. Finally, in view of the fact that $\varphi_\nu(\mathbf{x}) \propto \nu + \omega_0$, we obtain

$$Q \propto \left\langle \left| \sum_\nu g_\nu(\mathbf{x})\, g_{-\nu}(\mathbf{x}) \right|^2 \right\rangle.$$

In this expression $g_\nu(\mathbf{x})$ are obtained from $f_\nu^0(\mathbf{x})$ by a transformation of the type described by Eq. (12), which is unrelated to ν, so that the final dependence of $\overline{Q}$ on the width of the spectrum is the same as of the original quantity $Q^0 = |\sum_\nu f_\nu^0 f_{-\nu}^0|^2$, which describes conservation of the correlation observed in this case.

CONCLUSIONS

The above theoretical analysis shows that when certain conditions are obeyed, the conservation of the intraspectral correlation in the case of strong scattering by inhomogeneities can be explained in a simple manner. The loss of correlation of interest to us as a result of the scattering by static inhomogeneities is due to diffraction, which ensures that the scattering pattern is a function of the frequency, i.e., plays the role of the frequency dispersion of the investigated medium.

The observed conservation of the correlation could be explained by a weak diffraction-induced frequency dispersion of the scattering by small particles in the first case and by smooth irregularities in the second case. Obviously, the diffraction-induced frequency dispersion is most pronounced in the presence of narrow wave resonances in the scattering system, but there is no clear reason why this should happen in our experiments.

In explaining the experimental results we have ignored justification of the assumptions made that contribution of multiple scattering is negligible, the contribution of the interference in the fourth-rank field correlation functions, is unimportant, there is no long-range correlation in the case of rough surfaces, etc. However, in view of the fact that decorrelation is associated with all these factors and decorrelation was not observed, we assumed that such factors did not play a significant role in our experiments.

Stability of such an apparently "fragile" property as the correlation of the harmonics in the case of strong diffuse scattering is surprising. Therefore, it is appropriate to men-

tion a general cause of stability of the correlation found experimentally. The point is this: the correlation function of interest to us has the special form

$$Q \sim \sum_{\nu,\mu} U_\nu(\mathbf{r})\, U_{-\nu}(\mathbf{r})\, U_\mu^*(\mathbf{r})\, U_{-\mu}^*(\mathbf{r})$$

$$\times \exp\{i\varphi_\nu(\mathbf{r}) + i\varphi_{-\nu}(\mathbf{r}) - i\varphi_\mu(\mathbf{r}) - i\varphi_{-\mu}(\mathbf{r})\}.$$

Its reduction as a result of the scattering is due to dispersion of the phases as $\mathbf{r}$ increases. Clearly, the rate of this reduction is related also to the dependence of $\varphi_\nu(\mathbf{r})$ on ν. This factor changes the sum of the phases not in the first, but in the second order of the expansion in ν, i.e., the net phase varies as $(\nu-\mu)^2(\partial^2\varphi_\nu/\partial\nu^2)_{\nu=0}$, which may be a small quantity.

We conclude by noting that recording of higher correlation functions of the field has been used successfully in light scattering spectroscopy. The conservation of the spectral correlation in the case of strong diffuse scattering of light observed in the present study is an interesting effect also from the point of view of its use in spectroscopy of higher correlation functions,[6] which can be applied to study the specific processes destroying this correlation in the case of static and dynamic light scattering.

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