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Quantum Phase Noise and Field Correlation in Single Frequency Semiconductor Laser Systems

PHILIPPE B. GALLION, MEMBER, IEEE, AND GUY DEBARGE

Abstract—**The influence of quantum phase fluctuations which affect single frequency semiconductor lasers in various coherent detection systems is discussed in terms of photocurrent autocorrelation and spectral density functions. The general treatment given in this paper can be applied in diverse practical cases and points out the problems of phase correlation and phase matching between the two mixed optical beams. In the more general case the photocurrent spectrum is found to be composed of discrete and quasi-Lorentzian parts whose energies and spectral spreads are discussed as a function of the laser line width, the phase matching and the phase correlation between the two coherently combined fields.**

Introduction

THE temporal coherence properties of multimode semiconductor lasers have been discussed by a number of authors [1]-[8]. Experimental investigations of these properties were commonly achieved by using interferometric or beating techniques [9]-[11]. The higher temporal coherence of single longitudinal mode semiconductor lasers has also been discussed [3], [12]-[14] and experimentally investigated by the above mentioned methods [15]-[26]. For a single-frequency constant-amplitude laser source operating for above threshold, the random quantum phase fluctuations are usually considered to be the major source of optical spectral spread [3], [12], [17]-[19]. The phase noise of a single frequency semiconductor laser is a parameter of major importance in applications in which the temporal coherence properties of optical signals are involved. Such applications are transmissions systems and sensors in which phase or frequency modulated signals are detected by heterodyne or homodyne techniques [27]-[31] or discriminated by direct interferometric detection [32], [33]. The laser phase is also a determining parameter for incoherent systems in which a spurious coherent mixing produces a phase to intensity noise conversion: this is for example a case for the square law detection of a beam which is first split by a polarization dispersive medium, and the one part is later mixed through polarization sensitive elements, with the another undergoing a time delay [21].

The repercussion of spectral spread on such systems can be completely discussed in terms of the autocorrelation function and spectral density of the photocurrent; the purpose of this paper is to give a simple unified treatment which can be applied in the various above mentioned particular cases. The results are discussed as a function of the laser linewidth, phase matching, the balance and the phase correlation between the two mixed beams. In addition the phase noise sensitivity of various detection schemes is compared.

Manuscript received July 18, 1983; revised November 14, 1983.
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Quantum Phase Noise Model for Single Longitudinal Mode Semiconductor Lasers

The optical field emitted far above threshold by a biased single frequency laser is commonly modeled as a quasimonochromatic amplitude-stabilized field undergoing a phase fluctuation [12]

$$E(t) = E_o \exp j[\omega_o t + \phi(t)] \tag{1}$$

where ω_o is the average optical frequency and $\phi(t)$ is a stochastic process representing the random phase fluctuation leading to the broadening of the spectral line.

The time dependence of E_o on time t is an amplitude noise. Its contribution to the field spectrum can be neglected in spite of its large spectrum width because its integration power is much smaller than the phase noise one [12].

It is useful to introduce the optical field autocorrelation function defined as

$$\begin{aligned} G_E^{(1)}(\tau) &= \langle E^*(t) E(t+\tau) \rangle \\ &= \langle \exp [j\Delta\phi(t,\tau)] \rangle \exp (j\omega_o \tau) \end{aligned} \tag{2}$$

where $\Delta\phi(t,\tau)$ is the phase jitter, i.e., the random phase change between times t and $t+\tau$

$$\Delta\phi(t,\tau) = \phi(t+\tau) - \phi(t). \tag{3}$$

Under the above mentioned assumptions, the phase jitter $\Delta\phi(t,\tau)$ is usually assumed to be a zero-mean stationary random Gaussian process with the associated probability density function being [15]-[34]

$$W[\Delta\phi(\tau)] = \frac{1}{[2\P \langle \Delta\phi^2(\tau) \rangle]^{1/2}} \cdot \exp \left[- \frac{\Delta\phi^2(\tau)}{2 \langle \Delta\phi^2(\tau) \rangle} \right]. \tag{4}$$

$\langle \Delta\phi^2(\tau) \rangle$ is the mean-square phase jitter which is related to the instantaneous angular frequency-fluctuation spectrum $S_\phi(\omega)$ by [34]

$$\langle \Delta\phi^2(\tau) \rangle = \frac{\tau^2}{2\P} \int_{-\infty}^{+\infty} \left[\frac{\sin \frac{\omega\tau}{2}}{\frac{\omega\tau}{2}} \right] S_\phi(\omega) \, d\omega. \tag{5}$$

Using the well known relation [34]

$$\langle \exp [\pm j\Delta\phi(t,\tau)] \rangle = \exp [-\tfrac{1}{2} \langle \Delta\phi^2(\tau) \rangle] \tag{6}$$

therefore the laser field correlation function is expressed as

$$G_E^{(1)}(\tau) = \exp [-\tfrac{1}{2} \langle \Delta\phi^2(\tau) \rangle] \exp (j\omega_o \tau). \tag{7}$$

The case considered here concerns quantum phase-fluctuation affecting a far above threshold biased semiconductor laser: then the instantaneous angular frequency-fluctuation spectrum $S_\phi(\omega)$ can be assumed to be flat [12]-[16] leading to the mean-square phase-jitter $\langle \Delta\phi^2(\tau) \rangle$ increasing linearly with the time delay

$$\langle \Delta\phi^2(\tau) \rangle = 2\gamma |\tau| \tag{8}$$

where 2γ is the angular full linewidth at half maximum (FWHM) of the Lorentzian laser field spectrum $S_E(\omega)$ obtained by the Fourier transform of (7)

$$S_E(\omega) = E_o^2 \frac{\gamma/\P}{\gamma^2 + (\omega - \omega_o)^2}. \tag{9}$$

2γ is given by the Schawlow-Townes formula including the line broadening excess factor introduced by Henry [14] to explain the experimental values reported by Fleming and Mooradian [13]:

$$2\gamma = v_g^2 \frac{h\nu g n_{sp} \alpha_m}{4P_o} (1 + \alpha^2). \tag{10}$$

where v_g is the group velocity of the light in the active medium, $h\nu$ the lasing photon energy, n_{sp} the spontaneous emission factor [3], α_m the loss of the mirror, P_o the output power per facet and α the linewidth enhancement factor introduced by Henry [14].

For channeled-substrate-planar (CSP)-structure Ga Al As devices Fleming and Mooradian [13] have reported experimental values for the product $2\gamma P_o$ of $2\P \times 110 \times 10^6$ rad $\cdot$ s^{-1} $\cdot$ mW.

Recently authors [35]-[38] have discussed a modified model taking into account the relaxation resonance effect. A Lorentzian lineshape with a second order correction is then carried out. This correction is not taken into account in this study which refers only to the first order model (assuming a flat frequency-fluctuation spectrum leading to a rigorously Lorentzian line shape).

Total Detected Field Modelization

The purpose of this paper is to express the spectral density of the photocurrent produced by homodyne and heterodyne detection of two fields as a function of their linewidth, their relative weight and their remaining phase correlation. For this purpose we assume that the detected total field $E_T(t)$ is a superposition of a laser field $E(t)$ as expressed by equation (1) with a time-delayed and frequency-shifted image of itself:

$$E_T(t) = E(t) + \alpha E(t + \tau_o) \exp j\Omega t \tag{11}$$

where α is a real factor which accounts for the amplitude ratio between two mixed fields. For coherent optical communication systems in which a weak signal is detected with a powerful local oscillator we have $\alpha \ll 1$. For a phase-modulated optical signal detection by the retardation method [32], [33] α is closely one. α can take any value in the case of noncoherent single-mode transmissions systems and unbalanced sensors and when the modes coupling in a single-mode fiber with polarization dispersion is involved [21].

- τ_o is the time-delay which also refers to the eventual remaining phase correlation between the two combined beams. For coherent communication systems, τ_o is the time delay in the optical local field generator [23], [27], [30]. For differential phase shift keying systems and for phase-modulated optical detection by the retardation method [32], [33] τ_o is a bit duration. Finally τ_o is the optical time delay for unbalanced interferometric systems and the dispersion time for the polarization-dispersed modes in single mode fiber mixing [21]. The optical mixing of two fields of equal linewidth and uncorrelated phases is obtained by $\tau_o = \infty$.
- Ω is the mean frequency difference between two mixed fields. Ω is the intermediary frequency in heterodyne communication systems and the acousto-optic frequency in heterodyne sensors. $\Omega = 0$ for homodyne communications systems, for fiber interferometers without Bragg-cell, for the polarization dispersed modes in single mode fiber mixing and for the phase modulated optical detection by retardation method. With

appropriate values of α, τ_o and Ω (1) for the total detected field one can easily consider the above mentioned four cases.

Because of the square law of optical detectors, these interferometric systems are able to convert the quantum phase noise of the laser field $E(t)$ into intensity noise and then give a spectral spread of photocurrent $I(t)$. For stationary fields, the autocorrelation function $R_I(\tau)$ of the photocurrent depends only on the intensity correlation function of the detected total field $G^{(2)}_{E_T}(\tau)$ [8]

$$R_I(\tau) = e\sigma\, G^{(2)}_{E_T}(0)\, \delta(\tau) + \sigma^2\, G^{(2)}_{E_T}(\tau) \tag{12}$$

where

- e is the electron charge,
- σ is the detector sensitivity,
- $\delta(\tau)$ is the Dirac function,
- $G^{(2)}_{E_T}(\tau)$ is the first order optical intensity correlation function,

i.e., the second order optical field correlation function defined by

$$G^{(2)}_{E_T}(\tau) = \langle E_T(t) E_T^*(t) E_T(t+\tau) E^*(t+\tau)\rangle \tag{13}$$

where the brackets denote a time average.

The first term in (12) is the usual shot-noise associated with the dc component of the photocurrent and the second account for this dc component and fluctuations arising from phase into intensity noise conversion.

Photocurrent Spectrum in the Homodyne ($\Omega = 0$) Detection Case

The autocorrelation function of the photocurrent is obtained by substituting (11) into (13) with $\Omega = 0$. This substitution gives rise to 16 terms. Using (1) for the laser field after a straight-forward calculation allow us to express $G^{(2)}_{E_T}(\tau)$ only as a function of the phase jitter of linear combination of $\phi(t)$ and $\phi(t+\tau)$:

$$\begin{aligned} G^{(2)}_{E_T}(\tau) = E_o^4\ \{&(1+\alpha^2)^2 + 2\alpha(1+\alpha^2)\exp j\omega_o\tau_o \\ &\cdot \langle \exp j\Delta\phi(t+\tau,\tau_o)\rangle + CC + \alpha^2 \exp 2j\omega_o\tau_o \\ &\cdot \langle \exp j\Delta\Psi(t,\tau_o)\rangle + CC + \alpha^2 \\ &\cdot \langle \exp j\Delta\Phi(t,\tau_o)\rangle + CC\} \end{aligned} \tag{14}$$

where

$$\begin{aligned} \Psi(t) &= \phi(t+\tau) + \phi(t) \\ \Phi(t) &= \phi(t+\tau) - \phi(t) \end{aligned} \tag{15}$$

and CC denotes the complex conjugate of the preceding term.

$G^{(2)}_{E_T}(\tau)$ can be expressed only as a function of $\langle\Delta\phi^2(\tau)\rangle$ by using (6)

$$\begin{aligned} G^{(2)}_{E_T}(\tau) = E_o^4 \Bigg\{&(1+\alpha^2)^2 + 4\alpha(1+\alpha^2)\cos(\omega_o\tau_o) \\ &\cdot \exp\left[-\frac{\langle\phi^2(\tau_o)\rangle}{2}\right] + 4\alpha\cos(2\omega_o\tau_o) \\ &\cdot \exp\left[\langle\Delta\phi^2(\tau)\rangle - \langle\Delta\phi^2(\tau_o)\rangle - \frac{\langle\Delta\phi^2(\tau+\tau_o)\rangle}{2}\right. \\ &\left. - \frac{\langle\Delta\phi^2(\tau-\tau_o)\rangle}{2}\right] + 2\alpha^2 \exp\left[\langle\Delta\phi^2(\tau_o)\rangle\right. \\ &\left.\left. - \langle\Delta\phi^2(\tau)\rangle + \frac{\langle\Delta\phi^2(\tau+\tau_o)\rangle}{2} + \frac{\langle\Delta\phi^2(\tau-\tau_o)\rangle}{2}\right]\right\} \end{aligned} \tag{16}$$

For laser phase noise $\Delta\phi^2(\tau)$ is given by (8), using the reduced time delay, defined as $\overline{\tau} = 2\gamma\tau$ and noting by θ the mean phase difference $\omega_o\tau_o$ between the two mixed fields $G^{(2)}_{E_T}(\overline{\tau})$ expressed as

$$\frac{G^{(2)}_{E_T}(\overline{\tau})}{E_o^4} = \begin{cases} \left[(1+\alpha^2) + 2\alpha\cos\theta\exp\left(-\frac{\overline{\tau}_o}{2}\right)\right]^2 \\ \quad \text{for any } \overline{\tau} \\ + 4\alpha^2\exp(-\overline{\tau}_o)\ \{\sin h(|\overline{\tau}| - \overline{\tau}_o) \\ + \cos^2\theta\ [1 - \exp(|\overline{\tau}| - \overline{\tau}_o)]\} \\ \quad \text{for } 0 < \overline{\tau} < \overline{\tau}_o. \end{cases} \tag{17}$$

By virtue of the Wiener-Khintchine theorem, the spectrum of the photocurrent is given by the Fourier transform of its autocorrelation function. Omitting the shot noise term and using the reduced angular frequency, defined as $\overline{\omega} = \omega/2\gamma$, the two-sided spectrum $S_I(\overline{\omega})$ is written as

$$\begin{aligned} \frac{S_I(\omega)}{\sigma^2 E_o^4} = &\left[1 + \alpha^2 + 2\alpha\cos\theta\exp\left(-\frac{\overline{\tau}_o}{2}\right)\right]^2 \delta(\overline{\omega}) \\ &+ 4\alpha^2\exp(-\overline{\tau}_o)\ \frac{1/¶}{1+\overline{\omega}^2}\ \Bigg\{ch\overline{\tau}_o - \cos\overline{\omega}\,\overline{\tau}_o \\ &+ \cos^2\theta\left[\cos\overline{\omega}\,\overline{\tau}_o - \frac{\sin\overline{\omega}\,\overline{\tau}_o}{\overline{\omega}} - \exp(-\overline{\tau}_o)\right]\Bigg\} \end{aligned} \tag{18}$$

where $\delta(\overline{\omega})$ is the delta function.

Let us consider one of particular cases in which a two beams interferometer with equal optical field intensities is tuned up for a maximum optical power output. Therefore we have $\alpha = 1$ and $\theta = 2\,k¶$ with k integer. Taking moreover the Gaussian probability density function assumption for the phase jitter (4), one can find the following expression derived by Armstrong [15]

$$\begin{aligned} \frac{S(\overline{\omega})}{4\sigma^2 E_o^4} = &\left[1 + \exp\left(-\frac{\overline{\tau}_o}{2}\right)\right]^2 \delta(\overline{\omega}) \\ &+ \exp(-\overline{\tau}_o)\ \frac{1/¶}{1+\overline{\omega}^2}\ \left[sh\,\overline{\tau}_o - \frac{\sin\overline{\omega}\,\overline{\tau}_o}{\overline{\omega}}\right] \end{aligned} \tag{19}$$

The approximately Lorentzian part of this spectrum is plotted in Fig. 1(a) and 1(b) for various values of the parameters $\overline{\tau}_o$, i.e., the time delay to coherence time ratio, and $\theta = \omega_o\tau_o$, i.e., the mean phase difference between the two mixed beams.

Photocurrent Spectrum in the Heterodyne ($\Omega \neq 0$) Detection Case

In this paragraph we perform a quite similar analysis as in the previous chapter. The substitution of (11) into (13) give

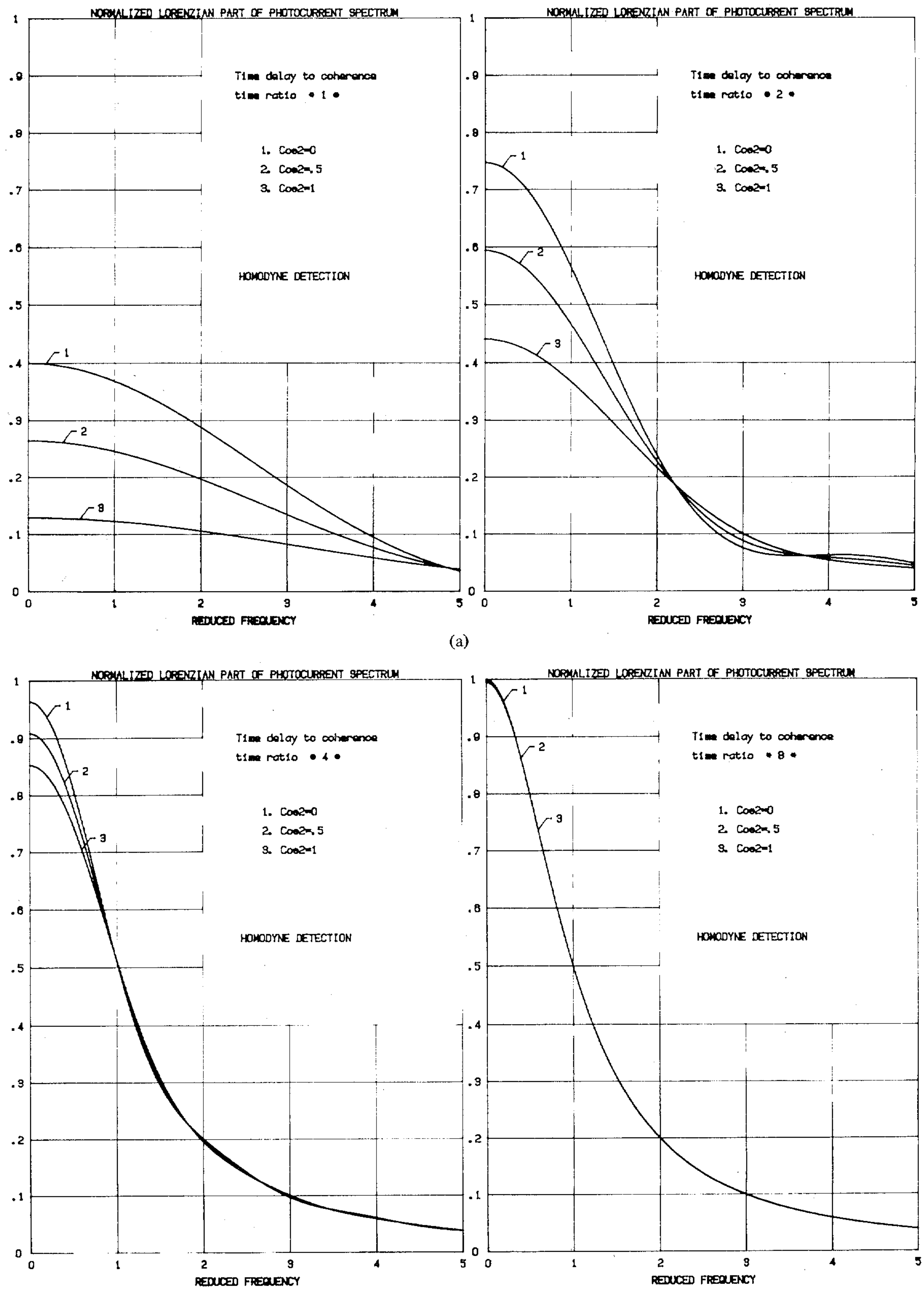

Fig. 1. Lorentzian part of the homodyne spectrum as a function (a) and (b) of the time delay to coherence time ratio $\overline{\tau}_o$ and of the phase matching parameter $\cos^2\theta$.

also rise to 16 terms but because an explicit time dependence appearing in the form of $\exp \pm j\Omega t$, 10 of them average out to 0; using the previous notations, the 6 surviving terms give the following result [22]-[23]:

$$G_{E_T}^{(2)}(\tau) = E_o^4 \{(1+\alpha^2)^2 + \alpha^2 \exp(j\Omega\tau) \cdot \langle \exp j\Delta\Phi(t, \tau_o)\rangle + \mathrm{CC}\} \tag{20}$$

where $\Phi(t)$ have the same meaning as above (15).

Using again the (6) $G_{E_T}^{(2)}(\tau)$ is written

$$G_{E_T}^{(2)}(\tau) = E_o^4 \left\{ (1+\alpha^2)^2 + 2\alpha^2 \cdot \cos \Omega\tau \cdot \exp\left[-\langle \Delta\phi^2(\tau_o)\rangle - \langle \Delta\phi^2(\tau)\rangle + \frac{\langle \Delta\phi^2(\tau - \tau_o)\rangle}{2} + \frac{\langle \Delta\phi^2(\tau+\tau_o)\rangle}{2} \right] \right\}. \tag{21}$$

In the same manner as the previous chapter, using the same reduced variables one obtains the first order optical intensity correlation function $G_{E_T}^{(2)}(\tau)$:

$$\frac{G_{E_T}^{(2)}(\overline{\tau})}{E_o^4} = (1+\alpha^2)^2 + \begin{cases} 2\alpha^2 \cdot \cos \overline{\Omega}\,\overline{\tau} \cdot \exp(-\overline{\tau}), & \text{for } |\overline{\tau}| < \overline{\tau}_o \\ 2\alpha^2 \cdot \cos \overline{\Omega}\,\overline{\tau} \cdot \exp(-\overline{\tau}_o), & \text{for } |\overline{\tau}| > \overline{\tau}_o \end{cases}. \tag{22}$$

Omitting again the usual shot noise term, the two-sided heterodyne photocurrent spectrum is then obtained by Fourier transform

$$\frac{S_I(\overline{\omega})}{\sigma^2 E_o^4} = [1+\alpha^2]^2\, \delta(\overline{\omega}) + \alpha^2 \exp(-\overline{\tau}_o)\, \delta(\overline{\omega} - \overline{\Omega}) + \alpha^2 \exp(-\overline{\tau}_o)\, \frac{1/¶}{1 + (\overline{\omega} - \overline{\Omega})^2} \cdot \left[\exp(\overline{\tau}_o) - \frac{\sin(\overline{\omega} - \overline{\Omega})\overline{\tau}_o}{\overline{\omega} - \overline{\Omega}} - \cos(\overline{\omega} - \overline{\Omega})\overline{\tau}_o \right]. \tag{23}$$

The approximately Lorentzian part of this spectrum is plotted in Fig. 2 for various values of the parameter $\overline{\tau}_o$, i.e., the time delay to coherence time ratio.

Discussion

The homodyne spectrum consists of two terms whose behavior relates closely to $\overline{\tau}_o$ and θ values. Let us examine the physical meaning of each term in several typical cases. The first term which is a dc component refers to the simple adding of the two optical powers and to the amount of remaining phase correlation between the two mixed beams. For large values of the normalized time-delay $\overline{\tau}_o$, i.e., for completely decorrelated fields, its dependence on the phase matching θ vanishes. On the other hand for a value of τ_o close to zero it becomes very sensitive on the θ value and it depends no longer on the spectral spread. The second term takes the form of an approximately Lorentzian lineshape which vanishes out for a close to zero $\overline{\tau}_o$ value. For large values of $\overline{\tau}_o$ this term stands for the optical mixing of two independent fields and it becomes rigorously Lorentzian with a FWHM which is twice the original laser linewidth. In this case the detector acts as an optical product-detector whose output is the autocorrelation product of the laser field spectrum. The spectrum is then no longer dependent on the phase matching θ.

The energies of the quasi-Lorentzian component $\epsilon_L(\overline{\tau}_o\theta)$ and of the dc component $\epsilon_\delta(\overline{\tau}_o\theta)$ of the spectrum are given by

$$\epsilon_L(\overline{\tau}_o\theta) = 2\alpha^2\, [1 - \exp(-\overline{\tau}_o)]\, [1 + \sin^2\theta \exp(-\overline{\tau}_o)] \tag{24}$$

$$\epsilon_\delta(\overline{\tau}_o\theta) = \left[1 + \alpha^2 + 2\alpha\cos\theta \exp\left(-\frac{\overline{\tau}_o}{2}\right) \right]^2. \tag{25}$$

For an in phase ($\theta = 0$) and an out of phase ($\theta = ¶$) optical mixing, the $\epsilon_L(\overline{\tau}_o, \theta)$ arrives at a minimum value. On the other hand in the case of a quadrature optical mixing, according to the well known fact that the phase noise conversion efficiency is optimum, the $\epsilon_L(\overline{\tau}_o, \theta)$ arrives at a maximum value.

The heterodyne case is quite different. Here the photocurrent spectrum consists of a dc component, a monochromatic and an approximately Lorentzian component centered at the intermediary frequency $\overline{\Omega}$. The first of them refers to the dc current associated with the noninteractive superposition of the two optical powers. The monochromatic one at the frequency $\overline{\Omega}$ expresses the remaining phase correlation between the two mixed beams. The third and last term is quasi-Lorentzian spectrum which cancels out for $\overline{\tau}_o = 0$ and become rigorously-Lorentzian with linewidth twice than the laser spectrum when $\overline{\tau}_o = \infty$.

Because of the mean frequency difference between the two-mixed beams the photocurrent spectrum is no longer dependent on the phase matching even in the correlated case. It is to be noticed that because of time averaging of phase matching in the heterodyne case the photocurrent spectrum is founded to be the same as in the homodyne case, with $\cos^2\theta = \frac{1}{2}$.

The energies in the quasi-Lorentzian part $\epsilon_{L,\Omega}(\overline{\tau}_o)$ and in the monochromatic part $\epsilon_{\delta,\Omega}(\overline{\tau}_o)$ of the spectrum are given by

$$\epsilon_{L,\Omega}(\overline{\tau}_o) = 2\alpha^2\, [1 - \exp(-\overline{\tau}_o)] \tag{26}$$

$$\epsilon_{\delta,\Omega}(\overline{\tau}_o) = 2\alpha^2 \exp(-\overline{\tau}_o). \tag{27}$$

These energies expressions stand for an uniform transfer of energy from the monochromatic to the Lorentzian component when increasing the decorrelation time $\overline{\tau}_o$.

As show in Fig. 3 there is no significant difference in the correlated spectral spread in the homodyne and heterodyne cases.

Some comparison of these results with experimental measurements has been made in various particular cases: Heterodyne detection of beams with uncorrelated phases [20], heterodyne detection of beams with correlated phases [23]-[24] and homodyne detection of beams with correlated phase and phase quadrature [21]. The results indicate for each case that the used theoretical description of the laser-phase random process is consistent with experiments. Experimental measurements in the more general case would be of a great interest.

Conclusion

The repercussion of quantum laser phase noise in various coherent systems has been discussed in terms of autocorrelation function and spectral density of the photocurrent. The given treatment includes the phase matching, the balance and the phase correlation between two coherently combined beams. For strongly correlated beams the photocurrent spectrum is found to be a Dirac function both in the homodyne and the heterodyne cases, and extremely dependent on phase matching

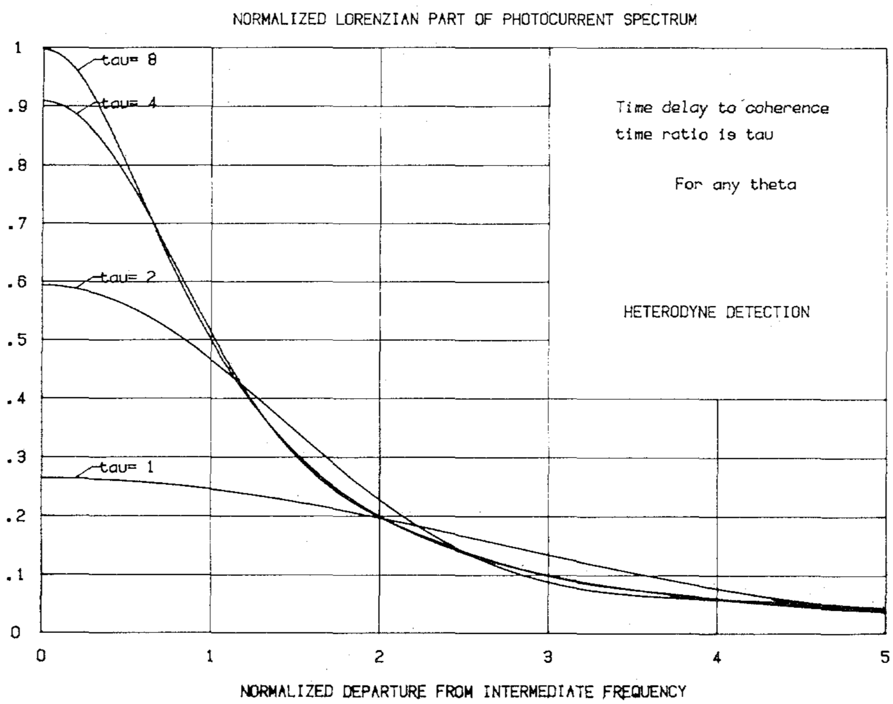

Fig. 2. Lorentzian part of the heterodyne spectrum as a function of the time delay to coherence time ratio $\bar{\tau}_O$.

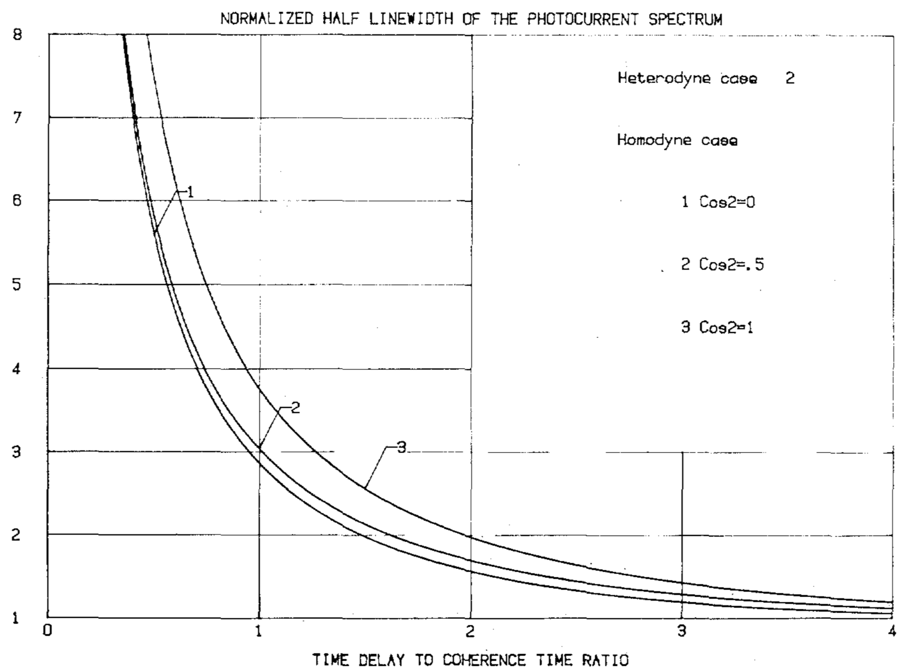

Fig. 3. Normalized half linewidth at half maximum of the Lorentzian part of the photocurrent spectrum.

in the homodyne case. For weakly correlated beams the spectrum is Lorentzian with a FWHM which is twice the original laser linewidth. Then in the homodyne case, the photocurrent spectrum depends no longer on the phase matching. The present analysis has been limited to laser sources exhibiting phase noise only. If the laser sources also exhibit intensity noise the photocurrent spectrum is a complex mixture of the two noises and its derivation seems to be quite tedious.

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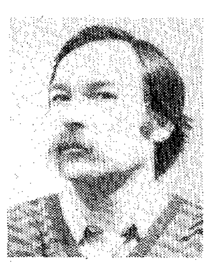

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