Source-Induced Noise in Optical Systems Driven by Low-Coherence Sources

MOSHE TUR, MEMBER, IEEE, EHUD SHAFIR, AND KJELL BLØTEKJÆER

Abstract—Recently available high-power quasimonochromatic light sources like light-emitting diodes, superluminescent diodes, and superfluorescent fibers may deliver enough power into optical and fiber-optic systems to make the source noise dominant at the system output. With an RMS value proportional to light intensity, this noise will limit the system sensitivity and dynamic range to levels which cannot be improved by injecting more light into the system. Source-related noise may originate not only from the source intensity fluctuations but also from a phase-to-intensity conversion process which is characteristic of many single-mode multiple-path optical systems. The shape of the power spectral density of the source-induced noise, being critically dependent on the physical structure of the system, is analyzed for a self-homodyne Mach–Zehnder structure and for a recirculating delay line. For single-path communication systems it is shown that source-originated noise exceeds both shot and thermal noise for a received optical power of only a few tens of microwatts.

I. INTRODUCTION

QUASIMONOCROMATIC, relatively wide-band optical sources have been recently proposed for single spatial mode optical systems which require light of very low temporal coherence. In the fiber-optic gyroscope, choosing a superluminescent diode (SLD) for the source minimizes drifts [1]. In fiber-optic sensor arrays utilizing the coherence-multiplexing technique [2], [3], the shorter the source coherence length, the higher is the achievable spatial density of sensors in the array. Many architectures used in fiber-optic signal processing require the use of low-coherence sources to ensure linear addition of intensities [4], [5]. Finally, the 10–100-μm coherence length of such wide-band sources can be very advantageous in quite a few measurement schemes [6], [7]. These developments also benefit from the recent achievements of high power in both SLD’s for sensor work [8], [9] and 1.3-μm LED’s for local area networks and short-haul communication systems [10]. Of great interest are also superfluorescent fiber sources which can emit as much as 10 mW of near infrared radiation in a single spatial mode [11]. However, all these sources are also noisy and their high-level optical output may change the noise balance in the system, shifting the dominance from thermal and shot noise toward the source noise.

Theoretically, it is common to model the (complex) output optical field of a thermal source as a sum of a huge number of statistically independent, randomly emitted contributions which form a random walk in the complex plane with Gaussian statistics [12]. In the absence of a better model, we will use this Gaussian model to describe the emission of the above quasimonochromatic sources, although it is experimentally clear that the behavior of SLD’s may slightly deviate from the pure thermal model [13]–[15]. In such a thermal-like source one can envision the complex optical field to strongly fluctuate both in amplitude and phase. The amplitude fluctuations will result in intensity noise at the system output and may limit the overall performance when this type of noise exceeds the always present shot noise and thermal noise. The situation becomes even worse in optical systems with multiple paths where the source phase fluctuations may be interferometrically converted to intensity variations at the system output [16]–[19], further degrading the system performance. Some of the implications of the use of thermal-like sources in coherence multiplexing schemes are discussed in [20].

This paper discusses the magnitude and shape of the power spectral density of the noise that results from the use of thermal-like sources in optical systems with and without multiple paths. Since multiple-path systems are known to act as RF filters [4], having intensity transfer functions with notches at specific RF frequencies, it will be of interest to study the interplay between noise filtration and noise generation in these systems. The mathematical formulation of Section II is followed by a few applications to some generic systems, Section III.

II. MATHEMATICAL FORMULATION

When spectral dispersion can be neglected, the output field of an optical system with polarization-maintaining, interferometrically interacting multiple paths can be generally expressed as a superposition of delayed versions of the input field $E_n(t)$

$$E(t) = \sum_{n=1}^{N} F_n E_n(t - T_n)$$

(1)

where $F_n (n = 1, 2, \ldots, N)$ is the complex weighting factor of the field which has experienced a (group) delay of $T_n$ while propagating through the system’s nth route $(1 \leq n \leq N)$. While the no-dispersion assumption is certainly valid for most free space systems, it is also acceptable in fiber systems with bounded differential delays $T_m - T_n$ [20].
The average output intensity is
\[
\langle I \rangle = \langle E \cdot E^* \rangle = \sum_{n=1}^{N} \sum_{m=1}^{N} F_n F_n^* \Gamma(T_m - T_n)
\]  
(2)

where \( \langle \cdot \rangle \) denotes ensemble averaging and \( \Gamma(t) = \langle E_n(t') E_n^*(t') \rangle \) is the self-coherence function of the input light [12]. Evidently, contributions to \( \langle I \rangle \) which come from terms for which \( T_m - T_n \) is much larger than the source coherence time \( \tau_c \) will be negligible.

The random fluctuations of \( E(t) \), and thus also of \( I(t) \), will manifest themselves as additional noise at the detector output current. According to the Wiener-Khinchine theorem, the power spectral density of \( I(t) \) is the Fourier transform of its autocovariance function \( C_I(t_1, t_2) \), given by
\[
C_I(t_1, t_2) = \langle (I(t_1) - \langle I \rangle)(I(t_2) - \langle I \rangle) \rangle
\]
(3)

Using (1) and changing the order of summation and averaging yields
\[
C_I(t_1, t_2) = \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n'=-1}^{N} \sum_{m'=-1}^{N} F_n F_n^* F_m F_m^* \cdot \langle E_n(t_1 - T_n) E_m^*(t_1 - T_n) \rangle \\
\cdot \langle E_n(t_2 - T_n) E_m^*(t_2 - T_n) \rangle - \langle I \rangle^2.
\]  
(4)

For polarized thermal-like sources, \( E_{in}(t) \) can be assumed to obey circularly complex Gaussian statistics [12] so that the forth-order moments appearing in (4) can be factored into the appropriate combinations of second-order moments [12]:
\[
\langle E(t) E^*(t) \rangle = \langle E(t_1) E^*(t_2) \rangle = \langle E(t_1) E^*(t_2) \rangle = \langle E(t_1) E^*(t_2) \rangle + \langle E(t_1) E^*(t_2) \rangle
\]
yielding, after some simplification
\[
C_I(t_1, t_2) = \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n'=1}^{N} \sum_{m'=1}^{N} F_n F_{n'} F_m F_{m'} \cdot \Gamma([t_1 - t_2] - (T_n - T_m)) \\
\cdot \Gamma([t_2 - t_1] - (T_n - T_m)).
\]  
(5)

It is interesting to compare (5) with the corresponding expression for \( C_I(t_1, t_2) \) that has been obtained for laser emission of the form \( E_0 \exp[(\omega t + \phi(t))] \), i.e., line broadening due to phase noise but with no intensity noise [17, eq. (22)]. These two equations are basically identical with two exceptions. (a) Naturally, the self-coherence functions have different forms. (b) While there are no contributions to [17, eq. (22)] from terms with \( n = m \) and \( n = m \), it is shown below that these terms which do appear in (5) represent the source intensity noise after being filtered by the optical system [4].

In order to further simplify the expressions for \( \langle I \rangle \) and \( C_I(t_1, t_2) \), we will assume that all optical routes within the system are of substantially different lengths with respect to \( \tau_c \). (Evidently, this excludes systems using coherence multiplexing which have been treated in detail in [20]). Under this assumption \( \Gamma(T_m - T_n) = 0 \), unless \( m = n \), and (2) reduces to
\[
\langle I \rangle = \sum_{n=1}^{N} \sum_{m=1}^{N} F_n^2 \Gamma(0) = I_0 \sum_{n=1}^{N} |F_n|^2
\]  
(6)

where \( I_0 = \Gamma(0) \) is the overall source optical power. Also, the term \( \Gamma([t_1 - t_2] - (T_n - T_m)) \) will contribute to \( C_I(t_1, t_2) \) of (5) only if
\[
T_n - T_m = T_m - T_n
\]
and
\[
t_1 - t_2 = T_n - T_m
\]  
(7)

where the \( \sim \) is interpreted as \( \sim \) within the order of \( \tau_c \). Strictly speaking, the equality sign in (7) could be relaxed to \( \sim \). However, since \( \tau_c \) for a thermal-like source is very small indeed, it is almost impossible to satisfy \( T_n - T_m \sim T_m - T_n \) to within \( \tau_c \approx 0.3 \text{ ps} \) (i.e., to within a path difference of less than or equal to 100 \( \mu \text{m} \)) unless strict equality is guaranteed by the basic design.

Thus (5) reduces to
\[
C_I(t_1, t_2) = \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n'=1}^{N} \sum_{m'=1}^{N} F_n F_n^* F_m F_{m'} \cdot \Gamma([t_1 - t_2] - (T_n - T_m)) \\
\cdot \Gamma([t_2 - t_1] - (T_n - T_m))
\]  
(9)

where we have used the fact that the self-coherence function \( \Gamma \) is Hermitian. Once in the form of (9), \( C_I(t_1, t_2) \) can be rewritten as a convolution
\[
C_I(t_1, t_2) = C_0(t_1, t_2) \ast \Gamma(t_1 - t_2)
\]  
(10)

\[
C_0(t_1, t_2) = \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n'=1}^{N} \sum_{m'=1}^{N} F_n F_{n'} F_m F_{m'} \\
\cdot \delta([t_1 - t_2] - (T_n - T_m))
\]  
(11)

where \( \delta([t_1 - t_2] - (T_n - T_m)) \) is the Dirac \( \delta \) function centered around \( T_n - T_m \).

Finally, the power spectrum of the source-induced noise is readily obtained by Fourier transforming (10), resulting in the product of two terms:
\[
S(f) = S_S(f) \cdot S_m(f)
\]  
(12)

where
\[
S_S(f) = \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n'=1}^{N} \sum_{m'=1}^{N} F_n F_{n'} F_m F_{m'} \\
\cdot \exp[2\pi i(T_n - T_m) f]
\]  
(13)
and

\[ S_a(f) = \int_{-\infty}^{+\infty} |\Gamma(t)|^2 \exp\left[2\pi i f t\right] \, dt \]

\[ = \int_{-\infty}^{+\infty} P(\nu) P^*(\nu - f) \, d\nu. \]  \hspace{1cm} (14)

\( P(\nu) \), being the Fourier transform of \( \Gamma(t) \), is the power spectrum of the source input light, and the last equality is based on the Fourier transform autocorrelation theorem [12].

While the second term in (12) \( S_a(f) \) is only source dependent, the first term \( S_0(f) \), being determined by the system itself, plays the role of a high-order transfer function. When the detector directly faces the source, we have \( N = 1, F_n = \delta_{n1} \), and from (13) and (12), \( S_0(f) = 1 \) and \( S(f) = S_a(f) \). Thus \( S_a(f) \) is really the power spectral density of the source intensity fluctuations.

Hence not only have we managed to express the total power spectral density of the output light intensity as a product of two terms, one describing the source and the other the system (see (12)), but we can gain even further insight by rewriting \( S(f) \) as

\[ S(f) = \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n=1}^{N} \sum_{m=1}^{N} F_n F^*_m F_n F^*_m \]  \hspace{1cm} (17)

\[ \cdot \exp\left[2\pi i(T_n - T_m) f\right] \cdot S_a(f) \]

\[ + \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n=1}^{N} \sum_{m=1}^{N} F_n F^*_m F_n F^*_m \]  \hspace{1cm} (15)

\[ \cdot \exp\left[2\pi i(T_n - T_m) f\right] \cdot S_a(f) \]

\[ = \sum_{n=1}^{N} |F_n|^2 \exp\left[2\pi i T_a f\right] \cdot S_a(f) \]

\[ + \sum_{n=1}^{N} \sum_{m=1}^{N} \sum_{n=1}^{N} \sum_{m=1}^{N} F_n F^*_m F_n F^*_m \]  \hspace{1cm} (16)

\[ \cdot \exp\left[2\pi i(T_n - T_m) f\right] \cdot S_a(f). \]

where \( \nu \) is the source center optical frequency and \( \Delta \nu \) is its half-power bandwidth. It is clearly seen that for either spectrum the 3-dB point of \( S_m(f) \) lies around \( f = \Delta \nu \). Since the optical sources discussed here have characteristic linewidths of \( \Delta \lambda \approx 10-100 \) nm, corresponding to \( \Delta \nu \approx 1.8-18 \) THz (at 1.3 \( \mu \)m), \( S_m(f) \) may be considered practically flat with \( S_m(f) = S_m(0) \) for all electronic frequencies of interest:

\[ S_m(0) = |\Gamma(0)|^2 \int_{-\infty}^{+\infty} |\Gamma(t)|^2 \, dt \]

\[ = \int_{-\infty}^{+\infty} |\Gamma(t)|^2 \, dt \]

\[ = t_0^2 \tau_c \]

\( S_m(0) = \frac{0.32 I_0^2}{\Delta \nu} \) (Lorentzian)

\( S_m(0) = \frac{0.66 I_0^2}{\Delta \nu} \) (Gaussian).  \hspace{1cm} (18)

The definition of the coherence time \( \tau_c \), and its relationship to the spectral widths were taken from [12].

At the output of the detector, the total noise is the sum of the source noise, shot noise due to the discrete nature of the detection process, and thermal noise generated at the detector load resistor [22]. Since both the shot and thermal random processes are also white at all electronic (RF) frequencies of interest, the total power spectral density of the detector output current, when normalized by the square of the average current is practically flat, given
by

$$
\xi(f) = \frac{S_n(0)}{I_0} + 2e \left( \frac{1}{R_{D}I_0} + \frac{\xi_{dark}}{[R_{D}I_0]^2} \right)
+ \frac{4kT}{[R_{D}I_0]^2} \tau.
$$

(19)

Here $R_D$ is the detector responsivity, $e$ is the electron charge, $I_{dark}$ is the detector dark current, $k$ is the Boltzmann constant, $T$ is the absolute temperature, and $R_L$ is the effective load resistance. The first term in (19), when expressed in [Hz]$^{-1}$, is actually the relative intensity noise (RIN) of the source [23].

Fig. 1 shows the dependence of $\xi(0)$ on the average incident optical power $I_0$ for different values of source linewidths and two values of the effective load resistance $R_L$. Also shown are curves of $\xi(0) \cdot V_3 \cdot I_0$ for pure shot noise and shot + thermal noise. The 1.3-μm source was assumed to have a Lorentzian lineshape. For a detector responsivity of $R_D = 0.5$ A/W and for the relevant range of values of $I_0$, the dark current of a single mode p-i-n diode could be completely neglected. When $I_0$ is large enough, the source noise becomes dominant, and consequently $\xi(0)$ approaches a constant equal to the source RIN. Under these conditions the output signal-to-noise ratio can be improved by increasing the input optical power. On the other hand, at low incident powers, thermal noise dominates. It is interesting to note that when the detector load resistance is 1 kΩ, source noise becomes dominant at $I_0 = 14$ dBm (25 μW) for a $\Delta\nu = 50$-nm source, and at $I_0 = 10$ dBm (10 μW) for a 10-nm source. These optical powers are indeed available these days even from LED's pigtailed to single-mode optical fibers. For a 50-nm source with $I_0 = 1$ mW, the overall source noise exceeds the shot + thermal noise by 17 dB. From now on we neglect both thermal and shot noise.

Using (6) and (12) the normalized source-induced noise at the output of a general system can be expressed as

$$
\xi_s(f) = \frac{S_n(f)}{\langle I \rangle^2} = \frac{S_n(f)}{\left( \sum_{n=1}^{N} |F_n|^2 \right)^2} \frac{S_n(f)}{I_0^2} \xi_s(f) \cdot \text{RIN}(f).
$$

(20)

$\xi_s(f)$ solely depends on the system parameters and plays the role of a normalized fourth-order transfer function. For the case of a single-path system:

$$
\xi_s(f) = 1.
$$

(21)

More complex systems will present more complicated $\xi_s(f)$, which will exhibit values either lower than unity as a result of filtration of the source intensity noise, or higher than unity due to phase-to-intensity noise conversion.

B. The Mach–Zehnder Architecture

For the Mach–Zehnder architecture [16], there are only two optical paths (i.e., $N = 2$) with a single time delay $\tau$. We analyze here only the incoherent case, $\tau >> \tau_c$.

Here, $G(t_1, t_2)$ of (10) comprises three duplications of $\langle \Gamma(t) \rangle^2$, centered around $-\tau, 0$, and $\tau$. The resultant normalized fourth-order transfer function $\xi_s(f)$ is readily found to be

$$
\xi_s(f) = \frac{|F_1|^4 + |F_2|^4}{(|F_1|^2 + |F_2|^2)^2} \cos(2\pi \tau f) + \frac{|F_1|^2 |F_2|^2}{(|F_1|^2 + |F_2|^2)^2} \cos(2\pi \tau f) + \frac{1}{2}.
$$

(22)

The term in the square brackets, which originates from the first part of (16), denotes the filtering effect of the Mach–Zehnder configuration on the intensity noise of the source [16], [21]. Indeed this term is the only term describing $\xi_s(f)$ when the optical outputs of the two arms are summed on an intensity basis, e.g., when the two outputs are orthogonally polarized or when the two outputs radiate two distinct areas of the detector. It is, therefore, the interferometric conversion of the source fluctuating phase which gives rise to the excess intensity noise described by the second term in (22).

Fig. 2 illustrates the previous discussion for a Mach–Zehnder configuration with equal contributions from the two arms ($|F_1| = |F_2|$), where (22) reduces to

$$
\xi_s(f) = \left[ \frac{1}{2} + \frac{1}{2} \cos(2\pi \tau f) \right] + \frac{1}{2}.
$$

(23)

The full expression of (23) is shown in curve $a$ of the figure. In the absence of the phase-induced intensity noise, i.e., without the last half, curve $b$ simply shows the filtering effect of the Mach–Zehnder architecture. Curve $c$ shows the constant value which $\xi_s(f)$ assumes for the single-path system, (21).

C. A Recirculating Structure

A general incoherent recirculating optical structure [19] (like the fiber-optic recirculating delay line and the incoherent Fabry–Perot configuration) is characterized by an infinite number of paths (i.e., $N = \infty$) and a single delay
time \( \tau (\gg \tau_0) \). For such a structure the complex field weighting factors (the \( F_n \) of (1)), except maybe for the first one (\( F_1 \)), form a geometric series

\[
F_n = T \cdot R^{n-1}, \quad n = 2, 3, \ldots \quad (24)
\]

where \( T \) is a constant and \( R \) is the roundtrip multiplication factor, which includes real loss and acquired phase. Thus any recirculating system can be characterized by the three complex constants, \( F_1, T, \) and \( R \). Let us also define the roundtrip intensity loss \( U = |R|^2 \). The above constants for some typical recirculating systems are given in [19, eqs. (6)–(8)], in terms of the physical parameters of the systems. For the general system the average output intensity can be found by substituting (24) into (6):

\[
\langle I \rangle = \left[ A_0 \right]^{1/2} \cdot I_0 = \left[ |F_1|^2 + |T|^2 \frac{U}{1-U} \right] \cdot I_0
\]

which also defines the constant \( A_0 \).

The autocovariance function \( C_\tau(t_1, t_2) \) is found to comprise infinite duplications of \( |\Gamma(t)|^2 \) centered around \( M \cdot \tau (M = 0, \pm 1, \pm 2, \ldots) \), and (13) and (20) yield

\[
\xi_\tau(f) = 1 + 2 \frac{\kappa}{A_0} \frac{U \cos (2\pi f \tau) - U}{(1 + U^2 - 2U \cos (2\pi f \tau))}
\]

(26)

where

\[
\kappa = \left| T F_1^* + |T|^2 \frac{U}{1-U} \right|^2
\]

i.e., the power spectrum is periodic, with a period of \( 1/(\text{delay time}) \). Equation (26) is identical (up to normalization) to the one obtained for the case of a recirculating structure driven by a semiconductor laser in which only phase noise was considered and amplitude noise neglected (see [19, eq. (21)]). However, the expression for \( \kappa \) is different. While for the case of pure phase noise [19], \( \kappa \) could be either positive or negative, depending on the system configuration; here \( \kappa \) is always positive. Consequently, for a thermal-like source, \( \xi_\tau(f) \) is always characterized by a maxima at integer multiples of \( 1/\tau \), whereas either minima or maxima have been observed for laser-driven recirculating structures [18], [19].

As an example, we consider now in more detail the case of the fiber-optic recirculating delay line (RDL) [17]. Both the output and the recirculating waves are examined; see the insets in Figs. 3 and 4. Table I lists the expressions for \( F_1, T, \) and \( R \) for the two waves (using [19], a similar table can be constructed for the transmitting and reflecting Fabry–Perot interferometer). The linear operation of the coupler is described in matrix form by

\[
\begin{bmatrix}
E_3 \\
E_4
\end{bmatrix} = \delta_0 \begin{bmatrix}
A & B \\
C & D
\end{bmatrix} \begin{bmatrix}
E_1 \\
E_2
\end{bmatrix}
\]

(28)

and the different optical fields \( E_i, i = 1, 4 \), are shown in the insets of Figs. 3 and 4.

It is interesting to note that as the loss of the RDL approaches zero, \( \kappa \) for the output wave vanishes, resulting in a flat unity-valued \( \xi_\tau(f) \). The interpretation is very simple: the intensity noise of the thermal source is filtered by the RDL producing the usual maxima at \( n/\tau \) and minima at \( (1/2 + n)/\tau \) [16]. But the RDL also generates


### TABLE I

<table>
<thead>
<tr>
<th>Output Wave</th>
<th>Recirculating Wave</th>
</tr>
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<tbody>
<tr>
<td>$F_1$</td>
<td>$\delta_o C$</td>
</tr>
<tr>
<td>$T$</td>
<td>$\delta_o A$</td>
</tr>
<tr>
<td>$R$</td>
<td>$\delta_o B \exp[-\alpha_o L]$</td>
</tr>
</tbody>
</table>

**Notes:**
1. $\delta_o$ is the coupler loss.
2. $\alpha_o$ is the fiber loss per unit length.
3. $L$ is the loop length.

Phase-inducd intensity noise with *minima* at $n/\tau$ and *maxima* at $(1/2 + n)/\tau$ [16]. The fact that these two processes fully compensate for each other indicates that in some sense thermal-like sources comprise equal "amounts" of intensity noise and phase noise. Note that in the case of the Mach–Zehnder configuration with equal contributions from the two arms (see (23)), the averaged filtered source intensity noise again equals the phase-induced noise.

In Fig. 3 we show $\xi_k(f)$ (see (26)) at the output of a fiber-optic recirculating delay line (see inset) driven by a thermal-like source, for $\delta_0 = 1$, $\exp[-\alpha_0 L] = 0.8$ and for different values of the intensity coupling ratio, $|B|^2$. Since a finite loss is assumed, the phase-to-intensity noise conversion process is apparently less effective than the filtering process, and peaks, through very shallow ($\pm 0.3$ dB), can be observed at $f = n/\tau$. Much stronger peaks ($\pm 8$ dB) are generated when the two above mentioned processes join forces as in the case of the tapped RDL of Fig. 4, where both terms of (16) exhibit peaks at $n/\tau$ (see also [19] and [20]).

### IV. CONCLUSIONS

The random nature of the output field of a low-coherence thermal-like source contributes noise to the output power of any optical system it drives. We have shown that the source intensity noise becomes significant (i.e., above shot and detector thermal noise) for a few tens of microwatts of optical power falling on the detector. For a detector with a load resistance of 1 kΩ, illuminated by 1 mW of optical power from a 50-nm-wide source, the overall source noise exceeds the shot + thermal noise by 17 dB. Recently introduced high-power sources of this type, such as light-emitting and superluminescent diodes, as well as superfluorescent fiber sources, easily deliver enough optical power to meet or exceed the above mentioned example.

These findings are of considerable practical importance: Once this source noise becomes dominant, the resulting signal-to-noise ratio, approximately given by (source linewidth)/(system video bandwidth), in (18), cannot be further improved by injecting more light into the system. As a result, the use of such thermal-like sources must be ruled out in optical systems requiring very high sensitivity and/or very wide dynamic range (e.g., inertial-grade fiber optic gyroscopes).

The optical system in between the source and detector may filter this source intensity noise by virtue of its intensity transfer function. In addition, the system may interferometrically convert the source random phase fluctuations into excess intensity noise at its output. In this work a general single-spatial-mode optical system was treated, and its effects on the source noise were studied. A general expression for the power spectral density of the output intensity noise was given in terms of a product of the input source noise and a fourth-order transfer function which solely depends on the system structure. The formalism was applied to a Mach–Zehnder configuration and to a fiber-optic recirculating delay line. In the former, the contribution of the phase-to-intensity conversion has a flat spectrum, with a level which equals the average level of the spectrally shaped filtered source intensity noise. In the latter, both contributions are spectrally shaped, and the resultant power spectral density depends on the actual structure of the system: it has a few dB variations for a tapped RDL (which is also equivalent to a transmitting Fabry–Perot), but is practically flat at the output of the device (the equivalent of a reflecting Fabry–Perot). This spectral flatness strengthens our interpretation that in some sense thermal-like sources contain approximately equal "amounts" of phase noise and intensity noise.

**Note Added in Proof:** Recently it has been experimentally established [24] that $E_1^*$ superluminescent fibers exhibit thermal-like noise characteristics which follow the predictions of this paper.

### REFERENCES


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