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Noise and photoconductivity in single quantum well infrared photodetectors

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ABSTRACT

The noise spectral density of current fluctuations in single quantum well infrared photodetectors is calculated using Langevin approach. Fluctuations of the incident photon flux are taken into account. The dark current noise spectral density has a Lorentzian shape (with a pedestal) with characteristic frequency equal to inverse time of the QW recharging. This effect is due to the modulation of the injection current by the charge in the QW. The noise gain and the photocurrent gain are expressed in terms of basic transport parameters. The ratio of the noise gain to photocurrent gain is different from unity in general case.

Keywords: Single quantum well infrared photodetector, noise spectral density, Langevin method, noise gain, photocurrent gain, QW recharging time, modulation of injection

1. INTRODUCTION

The noise of the dark current and photocurrent is an important factor in operation of the Quantum Well Infrared Photodetectors (QWIPs).[1] For typical operating conditions, the main source of fluctuations in QWIPs is generation-recombination noise associated with the excitation of carriers from the QWs into continuum and their capture into the QWs. The study of noise in QWIPs is important for applications to obtain the devices with high detectivity, and it also provides an additional physical insight for these systems. The fluctuations are the natural sources of the transient excitation, therefore the QWIP response to this excitation provides the information on internal physical processes, and thus forms the basis for the noise spectroscopy.[2,3] Due to the numerous experimental and theoretical works on noise properties of QWIPs with multiple QWs (see, e.g., Refs.[1,4–9]), the overall understanding of the QWIP noise characteristics is satisfactory. Some issues, however, are still controversial, such as the relation between the noise gain and the photocurrent (optical) gain, the frequency dispersion of the noise spectral density, and the role of the injecting contact.

In this paper we consider QWIP with single QW (SQWIP). SQWIP is especially attractive theoretically because its simple structure allows an accurate self-consistent calculation of the electric field, injection current, and charge accumulation in the QW, and hence a better understanding of the noise properties. In addition, SQWIPs are intrinsically fast devices promising for CO_2-laser based high-speed applications, for which the noise is an important consideration.[10–12] Moreover, SQWIPs can have highest responsivity and detectivity as compared to QWIPs with multiple QWs, provided that efficient optical coupling scheme is implemented. Although SQWIPs have been studied by several research groups (see, e.g., Refs. [13–16]), we are aware of only one theoretical paper[17] dealing with the noise in SQWIPs. However, the result obtained in Ref. [17] requires a number of restrictive assumptions (large photocurrent gain, absence of the electron transport from the QW to emitter, etc.).

In this paper we extend recent theoretical studies of SQWIPs[18–20] to present a noise theory for SQWIPs. The frequency-dependent spectral density of current fluctuations is expressed in terms of the basic transport and injection parameters, and is applicable for SQWIPs with different design concepts.

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2. MODEL

The SQWIP under consideration contains a single-level QW separated by undoped barriers from heavily doped contacts (Fig. 1). Our model basically follows Refs. [18] and [20], however, we do not assume any particular shape of the emitter and collector barriers, so for example the injection current I_e has a general dependence on the electric field E_e in the emitter barrier. An important transport parameter of the model is the efficiency β ($\beta \leq 1$) being the probability for an injected electron to pass from emitter directly to collector, while $(1-\beta)$ is the probability of electron capture by QW (here the meaning of the QW capture probability is different from that in the case of drift electron transport in QWIPs with multiple QWs – see Ref. [9]). The electrons emitted from the QW are collected by the collector and emitter with probabilities ζ and $(1-\zeta)$, respectively (electron transport from the QW to emitter is especially important for SQWIPs with triangular barriers[18–20] – see Fig. 1(b)). Electron transport across the emitter and collector barriers is assumed to be instantaneous, therefore, our analysis is limited by frequencies $\omega \ll v_T/\max(W_e, W_c) \sim 10^{12}\ \mathrm{s}^{-1}$, where v_T is a typical (thermal) velocity and W_e, W_c are the emitter and collector barrier thicknesses, respectively (the model of instantaneous jumps can be used for both thermoactivated emission and tunneling). We also neglect the electron interactions during traveling in the barrier regions, and single-electron

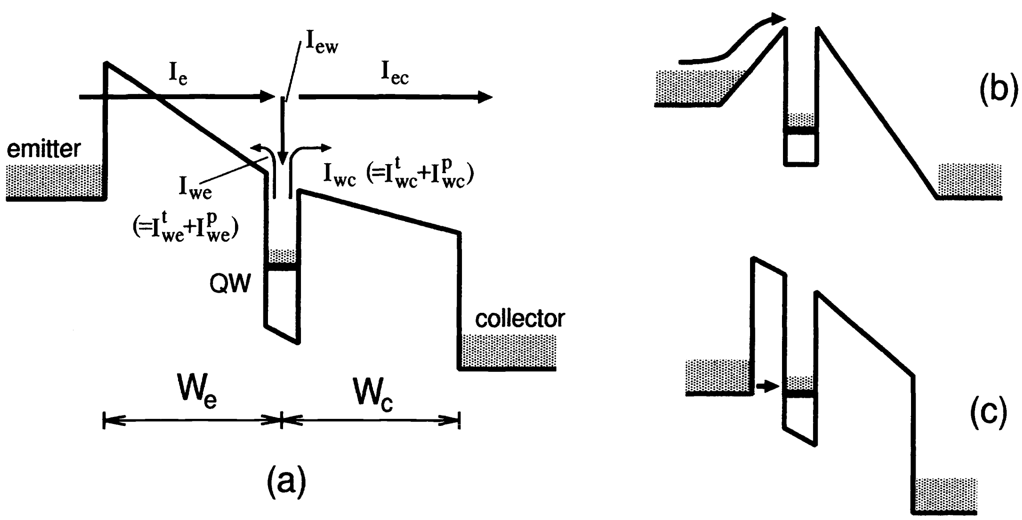

Figure 1. Schematic diagram of the conduction band profile and currents in an SQWIP with (a) thermionically assisted tunneling injection (rectangular barrier design), (b) thermionic injection (triangular barrier design), and (c) tunneling injection into QW (thin emitter barrier).

correlations. The dynamics of the electron transport in SQWIP is described by the following Langevin equations:

$$I = \frac{W_e}{W}\left(I_{ew} - I^t_{we} - I^p_{we} + \xi_{ew} - \xi^t_{we} - \xi^p_{we}\right) + \frac{W_c}{W}\left(I^t_{wc} + I^p_{wc} + \xi^t_{wc} + \xi^p_{wc}\right) + I_{ec} + \xi_{ec}\,, \tag{1}$$

$$\dot{Q} = I_{ew} - I^t_{we} - I^p_{we} - I^t_{wc} - I^p_{wc} + \xi_{ew} - \xi^t_{we} - \xi^p_{we} - \xi^t_{wc} - \xi^p_{wc}\,. \tag{2}$$

Here I is the total current through SQWIP (which is equal to the sum of the conduction and displacement currents at any cross section), $I_{ec} = \beta I_e$ and $I_{ew} = (1-\beta) I_e$ are the currents from the emitter to collector and well, respectively, and $I_{wc} = \zeta I_w$ and $I_{we} = (1-\zeta) I_w$ describe the electron transport from the well to collector and emitter

($I_w = I_w^p + I_w^t$ where superscripts t and p denote the currents due to thermo- and photoexcitation, respectively). The sign convention used in this paper is most simply understood if the electron charge is considered to be positive (this obviously does not affect the final results), then the direction of the current coincides with the direction of electron transport. Alternatively, one can consider true negative electron charge, then all currents (I_e, I_w, etc.) are also negative and $e < 0$. We neglect the current from the collector assuming sufficiently large bias voltage V. The well is assumed to be narrow, so the total structure thickness is $W = W_e + W_c$, while the finite well thickness in the first approximation can be taken into account via effective values for W_e and W_c. The current I_e depends mainly on the electric field E_e in the emitter, $E_e = E_e^0 + V/W - QW_c/(\varepsilon\varepsilon_0 AW)$, while I_w depends on the electron population in QW and electric fields in the barriers,[14,15,18] so all currents are some functions of the QW charge Q (here A is the SQWIP area, $\varepsilon\varepsilon_0$ is the dielectric constant, and E_e^0 is the parameter of SQWIP design). For simplicity we neglect the dependencies of β and ζ on the accumulated charge.

Studies of the steady-state characteristics and admittance of SQWIPs described by similar models have been reported recently.[14,15,18,20] In this paper we concentrate on the fluctuations which in Eqs. (1)–(2) are caused by random Langevin terms $\xi(t)$. All random terms (except ξ_{we}^p and ξ_{wc}^p – see below) have no mutual correlation, all of them are δ-correlated in time, and the corresponding spectral densities $S(\omega)$ are given by usual Schottky formula:

$$S_{\xi_{ew}}(\omega) = 2\,e\,\langle I_{ew}\rangle,\ S_{\xi_{wc}^t}(\omega) = 2\,e\,\langle I_{wc}^t\rangle,\ \text{etc.} \tag{3}$$

(brackets denote time averaging). The fluctuations of the photocurrent depend on the photon source noise. Let the photon flux incident to the QW has the spectral density $S_p(\omega) = 2\nu[1+\alpha(\omega)]$ where ν is the average flux (photons per second) and α describes the deviation from Schottky level. Then the noise of the photoexcitation currents from the QW is given by

$$\begin{aligned} S_{\xi_{wc}^p}(\omega) &= 2\,e\,\zeta\,\langle I_w^p\rangle[1+\zeta\,\eta\,\alpha(\omega)],\ \langle I_w^p\rangle = e\,\eta\,\nu, \\ S_{\xi_{we}^p}(\omega) &= 2\,e\,(1-\zeta)\langle I_w^p\rangle[1+(1-\zeta)\,\eta\,\alpha(\omega)], \\ S_{\xi_{we}^p\xi_{wc}^p}(\omega) &= 2\,e\,\langle I_w^p\rangle\,\zeta\,(1-\zeta)\,\eta\,\alpha(\omega) \end{aligned} \tag{4}$$

Here the last equation describes the mutual spectral density, and the absorption quantum efficiency η includes the finite probability for a photoexcited electron to escape from the QW. Equations (4) can be derived separating in the correlation functions the terms corresponding to one and two excitation events. The first term is proportional to the probability $\xi\eta$ or $(1-\xi)\eta$ while the second term is proportional to the corresponding product of probabilities. Equations (4) show that the photoexcitation current noise has the simple Schottky behavior only if the photon flux is Poissonian ($\alpha = 0$) or η is small.

3. RESULTS

Equations (1)–(4) allow us to calculate the noise properties of SQWIP. Applying the standard Langevin method[2,3] to the linearized version of Eqs. (1)–(2), we first formally solve Eq. (2) in the frequency representation taking into account the dependence of currents on the accumulated charge. Substituting the result into Eq. (1) and using Eqs. (3)–(4) for Langevin sources we obtain the following spectral density of the total current

$$\begin{aligned} S_I(\omega) = 2e\langle I_{ec}\rangle \, &+ \, 2e[\langle I_{ew}\rangle + \langle I_{we}^t\rangle \\ +\langle I_{we}^p\rangle(1+(1-\zeta)\eta\alpha(\omega))] &\left|\frac{W_e}{W} - \frac{\chi}{1-i\omega\tau}\right|^2 \\ +2e\left[\langle I_{wc}^t\rangle + \langle I_{wc}^p\rangle(1+\zeta\eta\alpha(\omega))\right] &\left|\frac{W_c}{W} + \frac{\chi}{1-i\omega\tau}\right|^2 \\ +4e\left(\langle I_{we}^p\rangle + \langle I_{wc}^p\rangle\right)\zeta(1-\zeta)\,\eta\alpha(\omega) \\ \times\mathrm{Re}\left[\left(\frac{W_e}{W} - \frac{\chi}{1-i\omega\tau}\right)\left(\frac{W_c}{W} + \frac{\chi}{1+i\omega\tau}\right)\right], \end{aligned} \tag{5}$$

$$\tau^{-1} = -\frac{dI_{ew}}{dQ} + \frac{dI_{we}}{dQ} + \frac{dI_{wc}}{dQ}, \tag{6}$$

$$\chi = \tau \left[-\frac{dI_{ec}}{dQ} + \frac{W_e}{W} \left(\frac{dI_{we}}{dQ} - \frac{dI_{ew}}{dQ} \right) - \frac{W_c}{W} \frac{dI_{wc}}{dQ} \right]. \tag{7}$$

Here the derivatives dI_i/dQ take also into account the dependence via the electric field modulated by Q. (Obviously $\langle I_{ew} \rangle = \langle I_{we} \rangle + \langle I_{wc} \rangle$.) One can see that in the case $\alpha(\omega) = const$ the spectral density has a Lorentzian shape (with pedestal) with the characteristic frequency τ^{-1}. Equation (5) can be used to determine the noise gain which is traditionally defined as $g_n(\omega) \equiv S_I(\omega)/4e\langle I \rangle$.

Equations (1)–(2) without noise sources can be also used to calculate the photocurrent gain (the ratio between the variations of the total current and the photoexcitation current for small-signal harmonic infrared excitation):

$$g_p(\omega) \equiv \frac{\delta I(\omega)}{\delta I_w^p(\omega)} = \zeta - \frac{W_e}{W} + \frac{\chi}{1 - i\omega\tau}. \tag{8}$$

The frequency dependence of the photocurrent gain is obviously governed by the same time constant τ of the QW recharging.[20] (This time constant corresponds to the characteristic time of establishing equilibrium at the injecting contact in QWIPs with multiple QWs.[21])

To simplify the further analysis let us assume $|dI_{we}/dQ + dI_{wc}/dQ| \ll |dI_{ew}/dQ|$ (that is a typical experimental case). Then $\tau = (1-\beta)^{-1}(-dI_e/dQ)^{-1}$ and $\chi = \beta/(1-\beta) + W_e/W$. If we also assume $\alpha(\omega) I_w^p \ll I_w$ (so that we can neglect the non-Poissonian term of the photocurrent), then the noise gain is given by

$$g_n(\omega) = g_n(\infty) + \frac{g_n(0) - g_n(\infty)}{1 + (\omega\tau)^2}, \; g_n(0) = \frac{1}{2}\frac{1+\beta}{1-\beta},$$
$$g_n(\infty) = \frac{1}{2} + (1-\beta)\frac{W_e}{W}\frac{W_e/W - \zeta}{\beta + \zeta - \beta\zeta}. \tag{9}$$

Under the same assumptions the ratio between the noise gain and photocurrent gain at small frequencies is given by the expression

$$g_n(0)/g_p(0) = (1+\beta)/[2(\beta + \zeta - \beta\zeta)] \tag{10}$$

(in conventional photoconductors this ratio is close to unity[11] while in our case unity is realized only if $\beta \to 1$ or $\zeta = 1/2$), and the minimal detectable photon flux ν_{min} at low frequency is given by

$$\nu_{min} \equiv \frac{\sqrt{S_I(0)\,\Delta f}}{e\eta\, g_p(0)} = \frac{\sqrt{2eI\,\Delta f}}{e\eta}\frac{\sqrt{1-\beta^2}}{\beta + \zeta - \beta\zeta}, \tag{11}$$

where Δf is the bandwidth.

4. DISCUSSION

Figure 2 illustrates the frequency dependence of the noise gain, the absolute value of the photocurrent gain, and signal-to-noise power ratio $|\delta I(\omega)|^2/S_I(\omega)$ for SQWIP with typical parameters. The frequency dependence of these quantities is obviously governed by the same time constant $\tau \simeq [(1-\beta)\, W_c/(\varepsilon\varepsilon_0 W) \times d(I_e/A)/dE_e]^{-1}$,[20] which is the time of the QW recharging, or establishing equilibrium at the injecting contact. The value of τ is inversely proportional to the derivative of the injection current with respect to the emitter electric field, and thus depends strongly on applied voltage, temperature, and SQWIP design. In the case of illumination by infrared radiation, τ does not depend on radiation intensity if the dark current is much larger than the photocurrent, and decreases with increasing intensity in the opposite case. The values of τ for typical SQWIP structures and operating conditions[20] are in the range $\sim 10^{-9} - 10^{-3}$ s. Because of strong (typically exponential) dependence of I_e on E_e, τ starts to decrease with illumination intensity when illumination changes E_e considerably (crudely this occurs when the photocurrent becomes comparable or larger than the dark current). The dependence of τ on temperature is typically exponential,[20] $\tau \propto \exp(-kT/\varepsilon_a)$, where ε_a is the activation energy, which can be used for evaluation of the QW parameters from the measurements of the SQWIP noise or photocurrent characteristics at different frequencies and temperatures.

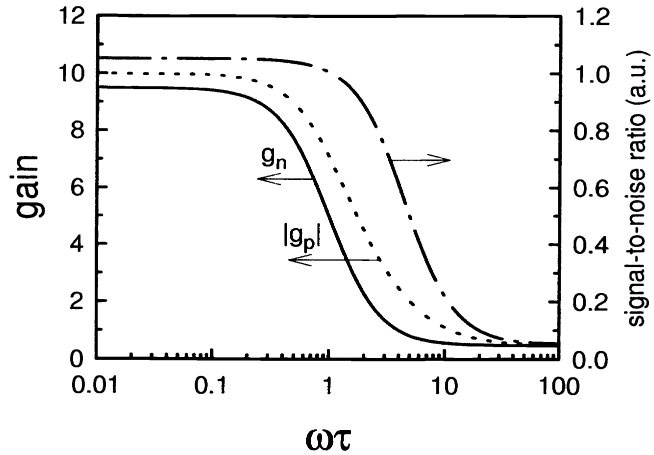

Figure 2. Illustration of the frequency dependence of the noise gain, photoconductive gain, and signal-to-noise power ratio for $\beta = 0.9$, $\zeta = 1$, and $W_e/W = W_c/W = 1/2$.

At high frequencies, $\omega \gg 1/\tau$, the QW recharging processes are "frozen", and the spectral density of current fluctuations is determined by the shot noise of elementary currents with appropriate geometrical factors. The photocurrent is due to the electrons emitted from the QW only, i.e. the primary photocurrent.[21] It is interesting to note that in the case $\zeta < W_e/W$ the high-frequency photocurrent gain $g_p(\infty)$ is negative.

At low frequencies, $\omega \ll 1/\tau$, the SQWIP operates in the quasistatic regime. The QW charge responds to the external excitation and modulates the injection current. In this regime the modulation effect results in a strong enhancement of both the photocurrent gain and the noise gain provided $1 - \beta \ll 1$.

If all the electrons injected from the emitter are captured by the QW ($\beta = 0$), then (see Eq. (9)) the low-frequency noise gain $g_n(0) = 1/2$ corresponds to the usual Schottky level. This has been observed experimentally in the SQWIP with thin emitter barrier (Fig. 1(c)).[16]

In the special case when $\zeta = 1$ and $\alpha = 0$ our main result given by Eq. (5) can be compared with the result of Ref. [17]. They coincide in the limit of high photocurrent gain, however, they are different for finite photocurrent gain because the result of Ref. [17] is not applicable in this case. We have checked that for $\zeta = 1$, $\alpha = 0$ the correct expression can be obtained also using the Fokker-Plank technique[3] (averaging $\exp(i\omega(t_m - t_n))$ where t_m and t_n are the moments of electron jumps). In the general case considered in the present paper, the Fokker-Plank technique becomes much more cumbersome than the Langevin method.

5. ACKNOWLEDGMENTS

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