Theory of asymmetric dual quantum well lasers

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A basic theory of asymmetric dual quantum well lasers is presented. The lasers consist of two quantum wells of different energy gaps, which are separated by a high and/or thick barrier layer. The barrier layer blocks carrier transport between the wells so that the rate of the transport becomes comparable to the rate of the radiative recombination. It is shown that this leads to novel phenomena of dual-wavelength lasing and wavelength switching with increasing injection current, in agreement with recent experimental results.

Laser diodes (LDs) with a variable emission wavelength $\lambda$ have wide applications. Conventional optical communication systems may require a tuning range $\Delta \lambda$ of less than 1% of $\lambda$, and LDs with a tunable distributed Bragg reflector are widely used. For other applications, such as future optical communication or optical recording systems, a wider $\Delta \lambda$ would be required. A solution is to use an external cavity, which enables us to get $\Delta \lambda/\lambda \sim 13\%$. However, the use of the external cavities places many restrictions on applications to small-size systems such as integrated optoelectronic systems. Hence, it seems important to develop monolithic LDs with a wide $\Delta \lambda$. As a new LD with this function, we previously proposed the asymmetric dual quantum well (ADQW) LD, for which we have experimentally demonstrated the discrete switchings of $\Delta \lambda = 13$ nm (Ref. 4) and 30 nm (Ref. 5), and quasi-continuous wavelength tuning of $\Delta \lambda = 22$ nm (Ref. 6), when $\lambda = 800$ nm. The purpose of the present letter is to present a basic theory of the ADQW LDs.

An ADQW LD consists of two quantum wells, well 1 and well 2, of different emission wavelengths $\lambda_1$ and $\lambda_2$ ($<\lambda_1$). The two wells are located in the core of a single optical waveguide as shown in Fig. 1, where well 2 of a wider band gap is located on the p-type side. If well 2 is located on the n-type side, roles of the electrons and holes should be interchanged in the following discussions. Let us make the thickness $d_2$ of well 2 larger than the mean-free path of energy relaxation processes of holes. Then, all holes injected from the p-type cladding layer first trapped by well 2, and they are then thermally activated to be transferred over the barrier into well 1. The point is to block this transfer by making the barrier high and/or thick, so that the rate of the transfer becomes comparable with the recombination rate. This should be contrasted with the case of normal multiple quantum well (MQW) LDs for which barrier layers are usually low and thin to inject carriers uniformly across many wells. If we used such a low and thin barrier in an ADQW LD, most injected carriers would occupy well 1 only, so that the carrier density $n_1$ of well 1 would not reach its threshold value at a reasonable value of $J$, resulting in sufficient optical gains both at $\lambda_1$ and $\lambda_2$.

As for electrons, well 1 of narrower energy gap is located on the n-type side. Hence, unlike the case of holes, the high and/or thick barrier layer tends to enhance the imbalance of the electron densities between the two wells. In order to supply well 2 with a sufficient density of electrons, we may make the potential at the edge of separate confinement layer on the n-type side (point A in Fig. 1) equal or higher than that of barrier layer, and also make the thickness $d_1$ of well 1 smaller than the mean-free path of energy relaxation processes of electrons. Then, a large part of electrons injected from the n-type cladding layer can go over well 1 into well 2, so that we can get sufficient electron densities in both wells. Another way of supplying electrons in well 2 is $n^+$ doping of the barrier layer. Once the sufficient electron densities are thus obtained in both wells, we can focus on holes and photons, for which we will write down rate equations. Although it is straightforward to extend the theory to include rate equations of electrons, we consider that main physics can be understood by the present simple equations.

Let $J_{21}$ be the hole-current density from well 2 to 1. Taking account of the general case when a part of $J$, $\xi J$, is directly injected into well 1, we may write the rate equations for $n_1$, $n_2$ and the photon densities $s_1^1$, $s_1^2$ of the lasing modes at $\lambda_1$, $\lambda_2$ as

$$\frac{dn_1}{dt} = \frac{\xi J + J_{21}(n_1,n_2)}{ed_1} - G_1^1(n_1)s_1^1 - G_1^2(n_1)s_1^2 = \frac{-n_1}{\tau_{n_1}}, \quad (1)$$

$$\frac{dn_2}{dt} = \frac{(1 - \xi)J - J_{21}(n_1,n_2)}{ed_2} - G_2^1(n_2)s_1^1 - G_2^2(n_2)s_1^2 = \frac{-n_2}{\tau_{n_2}}, \quad (2)$$

$$\frac{ds_1^1}{dt} = G_1^1(n_1)s_1^1 - \frac{s_1^1}{\tau_1} + \beta_1^1 \frac{n_1}{\tau_{n_1}}, \quad (3)$$

in turn increases. In other words, the quasi thermal equilibration for carrier distributions between the two wells is inhibited, which makes it possible to achieve $n_1 \sim n_2$ at a reasonable value of $J$, resulting in sufficient optical gains both at $\lambda_1$ and $\lambda_2$. 

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$$\frac{ds_1^1}{dt} = G_1^1(n_1)s_1^1 - \frac{s_1^1}{\tau_1} + \beta_1^1 \frac{n_1}{\tau_{n_1}}, \quad (3)$$
FIG. 1. Schematic band diagram of an ADQW LD. The shaded regions in the wells indicate the accumulated carriers.

\[
\frac{d\xi^2}{dt} = \left[ G_{1}^2(n_1) + G_{2}^2(n_2) \right] \beta_{1}^2 \frac{\rho_{1}}{\tau_{1}} + \beta_{2}^2 \frac{\rho_{2}}{\tau_{2}},
\]

(4)

where \( \tau_{1} \) and \( \tau_{2} \) are the recombination lifetimes in well 1 and 2, respectively, and \( \rho_{1} \) and \( \rho_{2} \) are the lifetimes of the \( \lambda_1 \) and \( \lambda_2 \) photons. The \( G_1 \), \( G_{2} \), \( \lambda_1 \), \( \lambda_2 \), \( \beta_{1} \), \( \beta_{2} \), \( \gamma_{1} \), and \( \gamma_{2} \) respectively denote the optical gain coefficients (the spontaneous emission factors) of \( \lambda_1 \) from well 1, \( \lambda_2 \) from well 1, and \( \lambda_2 \) from well 2. For example,

\[
\xi^2 = \frac{\gamma_{1} \rho_{1}^2}{\gamma_{2} \rho_{2}^2},
\]

(5)

where \( \Gamma_{1}^2 \) is the confinement factor of the \( \lambda_2 \) light with respect to well 1, \( c \) the light velocity, \( G_{1} \) the optical gain of well 1 at \( \lambda_2 \) and \( N_{1} \) the effective refractive index for the \( \lambda_2 \) lasing mode. The other factors, \( G_{1} \) and \( \beta_{1} \), do not enter the equations because well 2 has no quantum states at \( \lambda_1 \).

The above set of equations can describe both static and dynamic properties of ADQW LDs. To demonstrate that wavelength switchings indeed occur as a result of peculiar features of the above equations, we here present a solution for the static properties. Note that the equations do not take account of any effects of thermal heating, so that the wavelength switching is not due to the thermal heating. This means that the switching speed is not limited by slow times (~\( \mu \)s or longer) of thermal processes. The switching in the absence of the thermal heating has been experimentally confirmed by a recent time-resolved experiment.\(^5\)

To analyze the static properties, we set \( d/dt = 0 \) in Eqs. (1)–(4), and we may drop the spontaneous-emission terms in Eqs. (3) and (4). We also assume that \( \xi = 0 \). As for \( J_{21} \), although its rigorous functional form is unknown at present, we can say that \( J_{21} \) would depend only on \( n_2 \) when \( n_1 \) is not too large, because the thermal activation energy for \( n_1 \) is much larger than that for \( n_2 \). We also note that \( J_{21} \) is an increasing function of \( n_2 \). Wavelength switchings occur for any functional forms of \( J_{21} \) which satisfy these properties. For simplicity, we here assume the following form:

\[
J_{21} = \kappa_1 n_2,
\]

(6)

where \( \nu \) is the effective velocity of holes moving over the barrier. The barrier must be high and/or thick enough so that \( \nu \) becomes much smaller than that in normal MQW LDs. From Eq. (13) below, we find that \( \nu \) must be as small as ~\( 10^3 \) cm/s to achieve a reasonable value of \( J_{21} \).

By noting that \( G_{1}^1 \), \( G_{2}^2 \), and \( G_{2}^1 \) are also increasing functions of \( n_1 \) or \( n_2 \), we can easily find the following solution. Below the lasing threshold, \( s^2 = 0 \) and \( s^1 \) increases as

\[
s^1 = \frac{v}{\tau_{1}} \frac{(J - s^1)}{(J_{1}^1(n_{1th}) - G_{c})},
\]

(11)

As for the hole densities, \( n_1 = n_{1th} \) is fixed, whereas \( n_2 \) continues to increase as Eq. (8), until it reaches the threshold value \( n_{2th} \) given by

\[
G_{1}^2(n_{2th}) = 1/\tau_{1} - G_{c},
\]

(12)

at the second threshold current density

\[
J_{2th} = e(v + d_{2}/\tau_{2}) n_{2th},
\]

(13)

Here, \( G_{c} = G_{1}^2(n_{1th}) \) is the optical gain coefficient of well 1 at \( \lambda_2 \) (not \( \lambda_1 \)) for \( n_1 = n_{1th} \) which plays a crucial role in determining the switching characteristics, as shown below. When \( J < J_{2th} \) [see Eq. (16)], \( s^1 \) increases with \( J \) as

\[
s^1 = \frac{(J - J_{2th})}{(1/\tau_{2} - G_{1})},
\]

(14)

On the other hand, \( n_1 = n_{1th} \) and \( n_2 = n_{2th} \) are both fixed, and thus the hole supply to well 1, \( J_{21} \), is also fixed. This constant \( J_{21} \) should be balanced with the \( \lambda_1 \) emission plus the stimulated emission of the \( \lambda_2 \) photons from well 1 when \( G_{c} > 0 \) (or, when \( G_{c} < 0 \), minus the carrier generation in well 1 due to the absorption of the \( \lambda_2 \) photons). That is,

\[
\frac{v_{1}}{\tau_{1}} \frac{(J - J_{2th})}{(J_{2th} - G_{c})},
\]

(15)

Since \( s^1 \) increases with \( J \), \( s^1 \) in turn decreases when \( G_{c} = G_{1}^2(n_{1th}) > 0 \), which is a condition for the nonthermal switching of the emission wavelength. In fact, \( s^1 \) continues to decrease when \( G_{c} > 0 \) until it vanishes at \( J = J_{2th} \), where the third threshold current density

We here consider the case when the barrier is not too high so that the \( \lambda_1 \) light first starts lasing as \( J \) is increased.\(^4\)

The threshold value of \( J \) is found to be

\[
J_{1}^1 = e(v + d_{2}/\tau_{2}) n_{1th}/\nu n_{1th},
\]

(9)

where the threshold hole density \( n_{1th} \) is given by

\[
G_{1}^1(n_{1th}) = 1/\tau_{1}.
\]

(10)

When \( J_{1}^1 < J < J_{2th} \) [see Eq. (13)], \( s^2 = 0 \) and \( s^1 \) increases as

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s^1 = \frac{v_{1}}{\tau_{1}} \frac{(J - J_{1}^1)}{(J_{2th} - G_{c})},
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For current densities $J > J^2_{th}$, by expanding $n_1$, $n_2$, and $\delta^{21}$ around their values at $J = J^1_{th}$, we can easily see that $\delta^{21}$ continues to increase with $J$, and $n_1$ starts to decrease, whereas $n_2$ in turn starts to increase. The decrease of $n_1$ is because $n_1$ is consumed by the increasing stimulated emission of the $\lambda_2$ light [the third term in the right-hand side of Eq. (1)]. On the other hand, $n_2$ increases to compensate for the decrease of $n_1$ in order for the total optical gain at $\lambda_2$, $G^{21}_1 (n_1) + G^{22}_2 (n_2)$, to be kept constant. Since a small part of $J$ is consumed for this increase of $n_2$, the gradient $\partial n_1 / \partial J$ becomes slightly smaller than that for $J < J^1_{th}$. We have thus achieved a complete switching of the emission wavelength from $\lambda_1$ to $\lambda_2$, as schematically shown in Fig. 2.

In summary, we have presented the basic operation principles and the basic equations (1)–(4) of ADQW LDs. The high and/or thick barrier between two wells of different emission wavelengths blocks the carrier transport between the wells, resulting in dual-wavelength lasing and wavelength switching with increasing the injection current. Since this switching does not rely on any effects of thermal heating, the switching speed is not limited by slow times of thermal processes. The present theory agrees with recent experimental results.\(^4\)–\(^6\)