Beyond the absorption-limited nonlinear phase shift with microring resonators

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We show that the nonlinear phase shift produced by a ring resonator constructed from a given nonlinear optical material can be greater than the phase shift produced by a single pass through an infinite length of the same material when linear and nonlinear absorption are taken into consideration. The figure of merit (defined by the phase shift times the throughput) also improves for the ring resonator over that of the native nonlinear absorbing material. We finally show that these benefits of using the ring resonator as a nonlinear phase-shifting element can enhance the switching characteristics of a Mach–Zehnder interferometer.

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Nonlinear optical devices hold great promise for the implementation of dense, ultrafast all-optical signal processing. This promise has not yet been fulfilled, mainly because of materials limitations. Most optical switching devices operate by a π phase shift. Whereas phase shifts that are this large can be achieved with a variety of materials,[1] the corresponding activation energies and dissipation are too large to permit practical devices to be achieved. These two requirements are not complementary, as materials with large nonlinearity are typically highly absorbing, incurring loss and reducing phase shift.

One possible route to breaking this impasse is to use specialized device structures that provide artificial resonance enhancement. The nonlinear Fabry–Perot etalon, for example, has been studied extensively,[2] primarily as a bistable device. Here we study the nonlinear transfer function of ring resonators (RRs) acting as phase-shifting elements. These devices can be readily fabricated onto optical chips with moderately high finesse.[3] We define a figure of merit (FOM) that takes into account the nonlinear phase shift and loss and compare the RRs with nonlinear sections of various lengths. Over a range of incident intensities the achievable RR phase shift exceeds that of the nonlinear sections, even for infinite lengths, which produces the loss-limited maximum phase shift. In addition, for all intensities the FOM of the RR is higher. Finally, we use the RR and nonlinear sections as nonlinear elements in a Mach–Zehnder interferometer and show that the RR Mach–Zehnder interferometer can provide improved switching characteristics.

First we look at the nonlinear phase shift in the presence of linear and two-photon absorption. Writing the electric field as $E = A(z)\exp[j(k_0 z - \omega_0 t)]$, where $k_0 = \omega_0 n(\omega_0)/c$ is the linear phase constant and A is a complex amplitude, we obtain the following propagation equation, which describes the change in A as a function of distance[1]:

$$2jk_0 \frac{\partial A}{\partial z} + 2{k_0}^2 \frac{n_2}{n_0}(1 + jK)\,|A|^2 A + jk_0\alpha_0 A = 0\,. \quad (1)$$

Here n_2 is the nonlinear Kerr index, $K = \beta_2/2k_f n_2$, β_2 is the two-photon absorption coefficient, $k_f = 2\pi\nu_0/c$, and α represents absorption per unit length. This equation is valid for propagation in bulk but can be used for a channel waveguide with the appropriate spatially averaged quantities. Equation (1) has an exact solution, given by

$$a^2 = \frac{\exp(-\alpha L)}{1 + 2k_f n_2 K|E(0)|^2 L_{\mathrm{eff}}}\,, \quad (2a)$$

$$\phi = k_0 L + \frac{1}{2K}\ln[1 + 2k_f n_2 K|E(0)|^2 L_{\mathrm{eff}}]\,, \quad (2b)$$

where the input and output fields are related by $E(L) = a\exp(j\phi)E(0)$ and the effective interaction length that is due to linear absorption is $L_{\mathrm{eff}} = [1 - \exp(-\alpha L)]/\alpha$. There are three ways to increase the nonlinear phase shift: improve the material parameters (which may not be possible), increase the input intensity (which may not be practical), and increase the interaction length. If the interaction length $L \to \infty$, then $L_{\mathrm{eff}} \to 1/\alpha$, $I(L) \propto |E(L)|^2 \to 0$, and the nonlinear phase shift $\Delta\phi \to$ constant, thus bounding the phase shift.

Now, for a single-coupler RR,[4]

$$E(0) = \tau E(L) + j\sigma E_i\,, \quad (3a)$$

$$E_t = \tau E_i + j\sigma E(L)\,, \quad (3b)$$

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where E_i and E_t are the incident and the transmitted field amplitudes and τ and σ are the transmission and reflection coefficients at the coupler related by $\tau^2 + \sigma^2 = 1$, respectively, and the fields within the resonator are related by $E(L) = a \exp(j\phi)E(0)$ after one round trip. In the linear regime, the amplitude transfer function of the ring is

$$\frac{E_t}{E_i} = \frac{\tau - \exp(-\alpha L/2)\exp(jk_0 L)}{1 - \tau \exp(-\alpha L/2)\exp(jk_0 L)} \,. \tag{4}$$

Figure 1 shows the squared magnitude and phase of this transfer function, plotted versus the normalized detuning parameter $Q\delta\nu = Q(\nu_0 - \nu_m)/\nu_m$, where $\nu_m = mc/nL$ is the mth resonant frequency of the cavity and Q is the quality factor of the resonance. For all situations, the cavity circumference is $L = 50\ \mu\text{m}$ (which is typical of a compact, high-index ring), the ring absorption is $\alpha = 8\ \text{cm}^{-1}$ (assumed to be due only to a strongly absorbing material; ring scattering and radiation losses would be comparable but are neglected for simplicity), and the coupling coefficient is $\sigma = 0.2$, which results in slight overcoupling {critical coupling occurs at $\sigma = [1 - \exp(-\alpha L)]^{1/2} = 0.198$}. Operating near critical coupling maximizes the on–off contrast in nonlinear operation. The Q under these conditions is $\sim 1.2 \times 10^4$, and finesse $\mathcal{F}$ is ~ 78.

It was shown previously[4] that the nonlinearly induced phase shift scales as the square of the RR finesse for a single-coupled, lossless ring as a result of the cascading of two effects, intensity enhancement and cavity detuning. The intensity enhancement is due to the cavity buildup and generates a large, single-pass nonlinear phase shift. This phase shift detunes the cavity, causing a large change in the phase of the transmitted light. The maximum phase change depends on the initial detuning parameter; for negative detuning, the phase shift can approach 2π. For $\delta\nu = 0$, the maximum phase change is clamped at π, whereas, for positive detuning, the phase change is less than π. In the presence of linear absorption in the ring, a choice of coupling parameter that produces critical coupling or overcoupling results in the same phase-change properties as in the lossless ring, but cavity Q is decreased. To avoid bistability we use $\delta\nu = 0$, which results in a maximum π phase shift.

We define the FOM,

$$\text{FOM} = \left(\frac{\Delta\phi_t}{\pi}\right)\left(e\,\frac{|E_t|^2}{|E_i|^2}\right),$$

such that a π phase change over a $1/e$ intensity absorption length gives FOM $= 1$; clearly, FOM ≥ 1 is desired. Here, $\Delta\phi_t$ is the phase change of the transmitted light. In Fig. 2 we compare the FOM in the case of the RR ($\delta\nu = 0$) with nonlinear sections of lengths L, $10L$, and $\mathcal{F}L$ (which is the effective path length of the RR). For a given length, two-photon absorption limits the maximum value of the FOM for the nonlinear section, and, as the length goes to infinity, the FOM goes to zero. Two-photon absorption has a much weaker effect on the RR, which has a larger FOM over the entire range of incident intensities. Two-photon absorption figure $T = 8\pi K = 2.5$ exceeds the maximum desired value of unity[5] for a nonlinear material, so the nonlinear sections are expected to be strongly affected by two-photon absorption.

In Fig. 3 we compare the phase change and the throughput of the RR with nonlinear segments to better illustrate the improved FOM. At lower input intensities the throughput of the RR is zero, whereas the throughput of the nonlinear segments is limited by linear absorption. However, the RR begins detuning at low intensities, resulting in a phase change at the output; the nonlinear sections produce essentially no phase change. As the intensity increases, the RR throughput increases and the phase shift begins to saturate at the value determined by the initial detuning. The phase shifts of the nonlinear sections begin to increase, but their throughput decreases as a result of nonlinear absorption. The dotted curve in Fig. 3 shows the phase shift of a nonlinear section of infinite length, thus providing the maximum value allowable by absorption. Therefore, over a large range of input intensities, the RR can produce a greater phase shift than that which is normally allowed by absorption.

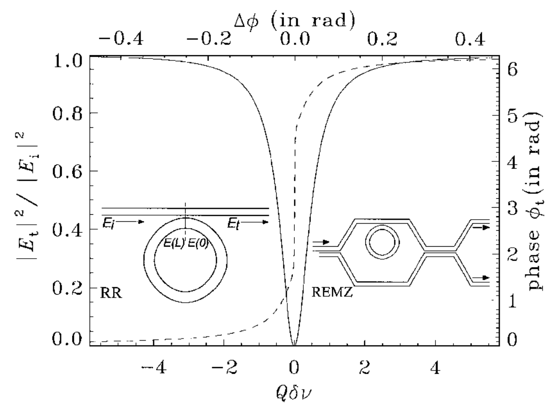

Fig. 1. Squared magnitude (solid curve) and phase (dashed curve) of the transfer function for a single-coupler RR near the critical coupling condition in which the rate of energy coupled into the cavity equals the rate of energy dissipated within the cavity. Insets, RR geometry used for Figs. 2 and 3 (left) and resonance-enhanced Mach–Zehnder interferometer geometry used for Fig. 4 (right).

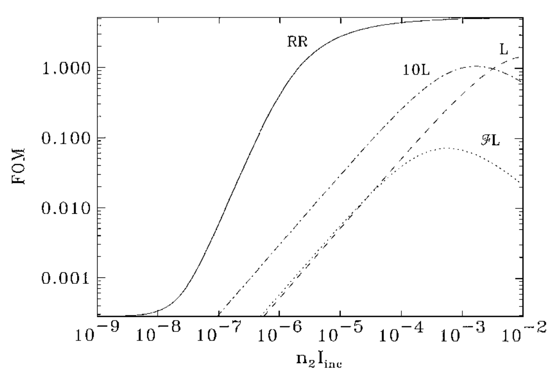

Fig. 2. Comparison of the FOM for the RR with $\delta\nu = 0$ (solid curve) and nonlinear sections of lengths L, $10L$, and $\mathcal{F}L$.

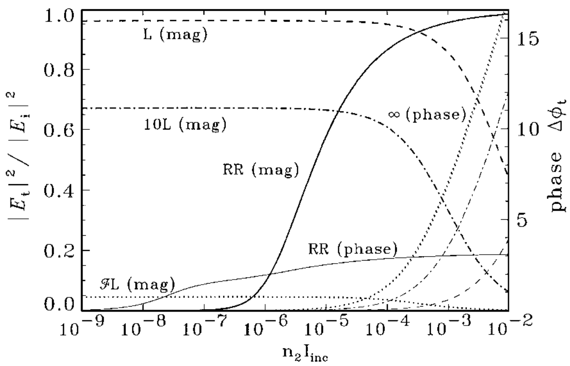

Fig. 3. Intensity transmission (heavier curves) for the RR (solid curves) with $\delta\nu = 0$ and nonlinear sections of lengths L, $10L$, and $\mathcal{F}L$. The corresponding phase changes are shown as thinner curves. The dotted curve corresponds to a nonlinear section of infinite length, which is essentially the same as for a length $\mathcal{F}L$.

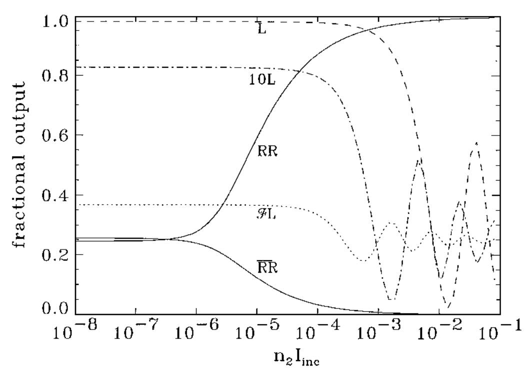

Fig. 4. Inverting and noninverting outputs of a Mach–Zehnder interferometer with a RR ($\delta\nu = 0$) used as a nonlinear element (solid curves) with inverting switching threshold $n_2 I_{inc} = 9 \times 10^{-6}$. Because of slight overcoupling, the outputs do not reach their expected values of 0.25 at low input levels. Inverting Mach–Zehnder outputs with nonlinear sections of lengths L (threshold, 6×10^{-3}), $10L$ (threshold, 7×10^{-4}), and $\mathcal{F}L$ (switching threshold not reached owing to large absorption) are shown.

The reason for the large phase shift of the RR in the presence of linear and two-photon absorption, and for the improvement in the FOM, is simple: The phase shift is the result of nonlinear refractive cavity detuning, which is not affected by linear absorption (which sets the cavity Q) and is weakly affected by nonlinear absorption. On resonance, at low input intensities, the incident light is highly dissipated within the cavity under critical coupling. A very small single-pass nonlinear phase shift within the ring (of the order of $2\pi/\mathcal{F}$, which can readily be achieved with even strongly absorbing media) can cause the RR to detune, resulting in an output phase shift that approaches π. Effects that are due solely to two-photon absorption are weak, as the RR FOM changes little with $K = 0$ compared with $K = 0.1$. In contrast, two-photon absorption has a much stronger effect on the nonlinear sections and is responsible for limiting the FOM as a function of intensity. Values of $K > 0.1$ do begin to affect the performance of the RR by reducing the FOM that is due to additional loss, but the RR degrades more slowly than does the nonlinear section.

Now, the question arises as to how one would take advantage of the phase-shift enhancement of the RR even though there is large loss at lower intensities (i.e., on resonance when $\delta\nu = 0$), over which the bulk of the phase shift occurs. One way to use the single-coupler RR is in a resonance-enhanced Mach–Zehnder device.[3,4] In this device the incident light is split into two arms of the interferometer, one of which contains the RR. Therefore, at low intensities, one fourth of the incident light exits each of the two outputs of the interferometer, as the light in one arm is completely dissipated within the cavity. At higher intensities, the inverting output drops to zero (while the other output rises) because of high transmission (matching the intensity in the other arm) through the RR with a saturating π phase shift. These characteristics are shown in Fig. 4, which also shows the balanced interferometer inverting output that results from nonlinear sections in one arm. Note that, without the RR, the output cannot go to zero (because of increasing loss with intensity owing to two-photon absorption, as opposed to decreasing loss for the RR owing to detuning) and is oscillatory because the phase shift can increase beyond π, whereas for the RR the maximum phase shift is clamped at π when $\delta\nu = 0$.

In conclusion, when it is viewed as a nonlinear phase-shifting element, a ring resonator can improve the achievable phase shift and figure of merit compared with that of the native material in the presence of linear and nonlinear absorption. These properties may allow practical switching and logic devices to be made from highly nonlinear materials that would otherwise be unsuitable.

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