Optical Nonreciprocity of Cold Atom Bragg Mirrors in Motion

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Reciprocity is fundamental to light transport and is a concept that holds also in rather complex systems. Yet, reciprocity can be switched off even in linear, isotropic, and passive media by setting the material structure into motion. In highly dispersive multilayers this leads to a fairly large forward-backward asymmetry in the pulse transmission. Moreover, in multilevel systems, this transport phenomenon can be all-optically enhanced. For atomic multilayer structures made of three-level cold $^{87}$Rb atoms, for instance, forward-backward transmission contrast around 95% can be obtained already at atomic speeds in the meter per second range. The scheme we illustrate may open up avenues for optical isolation that were not previously accessible.

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Much attention has been devoted to the development of advanced materials and composite systems to achieve optical functionalities not readily available in natural media. Such optical metamaterials can be engineered to stretch the rules that govern light propagation and light-matter interaction, potentially seeding a new paradigm in all-optical, optoelectronic, and optomechanical devices. Photonic crystals and negatively refracting media are prominent instances of manmade systems, the optical properties of which can be tailored to a great extent. Nevertheless, some tasks are more difficult than others. Already in the familiar process of linear reflection and transmission of light [1] it is in general hard to achieve nonreciprocity. In particular, multilayer photonic structures made with linear isotropic media with dissipation and/or gain may exhibit reflection nonreciprocity, i.e., an unbalance between the forward and backward reflectivities. There are even cases in which one of them can be made negligible and the other one can increase without limit with the sample thickness [2]; this occurs in the so-called $PT$-symmetric media that exhibit a variety of peculiar optical properties [3,4]. Yet, it is not possible to achieve a nonreciprocal transmissivity in such linear and passive systems. Transmission reciprocity is almost ubiquitous in optics [1,5].

Nonreciprocal transmission is, however, rather desirable for information processing and crucial to the development of optical-based functional components in photonics. In much the same way in which electrical nonreciprocity has been realized through diodes, devising an optical diode is challenging, even in theory. Ideally, an optical diode would allow total light transmission over a bundle of wavelengths in one direction, providing total isolation in the reverse direction. Previous work on nonreciprocal transmission has been based on either magneto-optical effects [6–8] or nonlinear processes [9,10]. Other mechanisms have also been explored [11–13] and realized experimentally [14], including more involved diode designs based on two-dimensional square-lattice photonic crystals [15] or non-symmetric photonic crystal gratings exhibiting anomalous diffraction effects [16]. Interesting proposals to achieve full nonreciprocity have recently been proposed either by dynamically inducing indirect photonic transitions [17] or by exploiting other acousto-optic effects induced by ultrahigh frequency coherent acoustic waves [18]. The latter two are instances in which the time independence of the optical response is broken due to the presence of the phonon field. However, most of these reported schemes to date require complex structures, demanding operating thresholds often with fairly low nonreciprocal output. Most importantly they exhibit a limited tuning range, which clearly restrains the efficiency of such schemes.

As dynamic optical isolation is in great demand in advanced optical communications, the question we will address here is whether a different physical principle can be used to devise a scheme that provides all-optical control over nonreciprocal transmission. An efficient way to create optical isolation is via time-reversal symmetry breaking [19,20]. Within such a regime several recent schemes have demonstrated good properties. However, these systems exhibit little room for tuning, and either large losses [5] or magnetic fields that could hamper the performance of nearby devices [21]. In contrast, we propose a mechanism where a large and tunable nonreciprocal transmission effect can be observed in a linear, dispersive, and moving Bragg multilayer. The motion of the multilayer breaks the
time-reversal symmetry. This is used to control nonreciprocity, which may be further enhanced by optically controlling the degree of dispersion of the multilayer. Bringing together motion and control over optical dispersion constitutes a new physical approach for nonreciprocal transmission management in multilayer photonic structures.

To start with, we consider light propagating at normal incidence on a planar inhomogeneous multilayer made of isotropic, linear, nonmagnetic materials with non-negligible losses, placed in an otherwise empty space and moving perpendicularly to the multilayer-vacuum interface (i.e., along the optical axis of the system). This provides the simplest example of forward-backward transmission asymmetry induced by macroscopic motion. While in the multilayer rest frame the frequency-dependent transmission $T(\omega)$ is reciprocal, in the lab frame at a given incident frequency $\omega'$, the transmission $\tilde{T}(\omega')$ is in general different in the two propagation directions since the Doppler effect gives rise to two different rest frame frequencies. Specifically, consider a multilayer moving with constant velocity $v = \hat{x}(v/c \ll 1)$ along its optical axis $x$ (see Fig. 1). The frequency $\omega'$ of a monochromatic light beam propagating along $\pm \hat{x}$ in the lab frame (primed) is seen in the rest frame (unprimed) as $\omega'_\pm = (1 \mp v/c)\omega'$ (first order in $v/c$). More generally for a Gaussian pulse of central frequency $\omega_p$ and spatial length $L'$ propagating forward ($+\hat{x}$) and backward ($-\hat{x}$) across the moving multilayer, the corresponding lab frame transmission $T_p$ can be written as

$$T_p(\omega_p, L') \approx \frac{L'}{c\sqrt{2\pi}} \int_{-\infty}^{\infty} d\omega' e^{-[(\omega'_- - \omega_p)^2/(2c/L')^2]}$$

$$\times T\left((1 \mp \frac{v}{c})\omega'_p\right). \quad (1)$$

Thus, forward-backward asymmetry in pulse transmission can be characterized by the nonreciprocity parameter $\Delta T_p(\omega'_p, L') \equiv T_+ - T_-$. A physical insight into the mechanism we propose may already be provided by considering a multilayer having a smooth frequency dependence of $T(\omega)$. Expanding $T(\omega)$ in a Taylor series one has

$$\Delta T_p(\omega'_p, L') \approx -\frac{v}{c}[2\omega'_p(dT(\omega'_p)/d\omega) + (2c/L')^2(d^2 T(\omega'_p)/d\omega^2)],$$

which shows that for a nearly monochromatic pulse ($L' \rightarrow \infty$) an appreciable degree of nonreciprocity would require a steep frequency derivative of $T(\omega)$. However, values not far from $|dT(\omega)/d\omega| \sim c/(\omega_p v)$ are hardly achievable via standard multilayer systems such as, e.g., solid state dielectric Bragg mirrors [24] where nonreciprocity values of $|\Delta T_p|$ are expected to be far from unity.

Such a small values of $|\Delta T_p|$ can be increased by using ultracold atoms in optical lattices [25], which have been the playground over the years for the study of complex optical phenomena [26,27]. In the following we will consider an atomic photonic structure (see Fig. 1) whose constituent materials can be arranged so as to form a periodic multilayer structure that effectively behaves as a Bragg mirror [28], and which can be set into motion at speeds of several meters per second [29,30]. Such an atomic Bragg mirror naturally exhibits resonant absorption and dispersion and represents an interesting challenge...
within the context of cold atoms optomechanics [31–33]. Such a system can be realized through cooling atoms into the vibrational ground state of each potential well of a sufficiently long and red-detuned confining 1D optical lattice (dipole trap). The resulting atomic distribution may be described as an array of disks separated by vacuum (see Fig. 1) whose thickness \( d \) is essentially given by the rms position spread around the minima of the optical potential and is much smaller than the lattice periodicity \( a \) [34]. The corresponding periodically modulated refractive index has been predicted to give rise to pronounced photonic stop bands [35–38]. In spite of the fact that efficient photonic structures with cold atoms are difficult to make because dilute atomic clouds have a low refractive index contrast and intrinsically short lengths, efficient Bragg reflection has recently been demonstrated in atomic multilayers that yield a well-developed one-dimensional photonic gap with nearly 80% reflectivity [28].

Such atomic Bragg mirrors are fairly robust against imperfections such as varying density across the mirror length and can be set to move making use of schemes that are familiar to experiments exploring the wave packet dynamics in light-shifted potentials [29,30]. The motion of the periodic atomic multilayers can in fact be envisaged as a collective oscillation of the atomic wave packet in the light-shift potential wells, much resembling that of a coherent state of atoms oscillating within a harmonic potential. Such oscillations may be induced, e.g., by suddenly shifting the optical lattice after the atoms have been brought to equilibrium at the bottom of the lattice sites [29]. Not only can such an atomic mirror lead to unidirectional transmission effects over 95%, but nonreciprocity can also be effectively tuned. Such a control (tuning) of the transmission nonreciprocity can be performed all optically through implementation of a three-level atomic configuration in the presence of a strong pump beam (see Fig. 1).

We refer, to be definite, to the D2 line of \(^{87}\text{Rb} \) atoms. The dressed atom susceptibility [27] is then characterized by an Autler-Townes (AT) doublet rather than by a single absorption line and can be tuned on demand by changing the pump intensity and frequency as described by

\[
\chi(\omega_p) = \frac{n \mu_{13}^2}{2 \varepsilon_0 \hbar} \frac{\gamma_{12}^R - i (\Delta_p - \Delta_e)}{(\gamma_{13}^R - i \Delta_p) + \Omega_e^2},
\]

where \( \gamma_{12} \) (\( \gamma_{13} \)) is the dephasing rate of the spin (optical) coherence \( \rho_{12} \) (\( \rho_{13} \)), \( \Delta_p = \omega_p - \omega_{31} \) (\( \Delta_e = \omega_e - \omega_{32} \)) is the probe (pump) detuning relative to transition \( |1\rangle \leftrightarrow |3\rangle \) \((|2\rangle \leftrightarrow |3\rangle)\), while \( \mu_{13} \) denotes the relevant electric-dipole moment. We assume a pump of Rabi frequency \( \Omega_p \) traveling in the \( y \) (or \( z \)) direction and a probe of Rabi frequency \( \Omega_p \) traveling in the \( x \) direction (the probe beam may make a small angle with the \( x \) direction so to adjust the photonic crystal periodicity as in [28,37]). The dressed susceptibility in Eq. (2) holds in the limit of a weak probe \((\Omega_p \ll \Omega_e)\), within the rotating-wave and electric-dipole approximations.

In particular, for \( \Omega_e \gg \gamma_{13} \), the dressed susceptibility evolves into that for two well-separated AT split levels. The dressed atoms are loaded into an optical lattice of period \( a \) formed by retroreflecting a light beam of wave-length \( \lambda_L = 2a \) (see Fig. 1). As mentioned above, in each period of such a dipole trap the confined atoms occupy a small region \( d \ll a \) with a homogeneous volume density \( n \).

This periodic 1D index modulation gives rise to pronounced photonic stop bands as expected, and the transfer matrix formalism can be used to describe the propagation of a probe field of frequency \( \omega \) through such an atomic stack of length \( L \gg a \), to compute the transmissivity \( T(\omega) \) in its rest frame (see [37] for a detailed description in an analogous context). Notice that in a nonresonant multilayer stack, as in the case of a standard distributed Bragg reflector, a single stop band opens up around the Bragg frequency; conversely, in the presence of a resonant absorption line close to the Bragg frequency, two stop bands will originate from the interplay between the polaritonic gap due to the resonance and the Bragg gap due to the periodic index modulation [37]. Yet, when the pump is on, a third narrow stop band arises and fits between two wider ones separated by the AT splitting [39,40]. Precisely this additional stop band is here exploited as its edge gives rise to a very steep frequency dependence of the transmissivity \( T(\omega) \) near the atomic resonance where absorption is significantly suppressed. The width and position of this narrow stop band depend critically on both pump frequency \( \omega_p \) and pump intensity \( \propto |\Omega_p|^2 \) [see Fig. 2(a)], which is what allows an efficient optical control of the degree of reciprocity in the transmissivity \( T(\omega) \). Figure 2(b) shows the nonreciprocity parameter \( \Delta T_p \) for a light pulse incident onto a moving dressed atom photonic crystal (see Fig. 1) as a function of the pulse detuning \( (\omega_p' - \omega_{31})/\gamma_{13} \) of central carrier frequency from atomic resonance. A very large nonreciprocity effect appears even at a speed as slow as \( v = 20 \text{ m/s} \) owing to the steep frequency dependence of the rest frame transmissivity \( T(\omega) \), while all optical control of \( \Delta T_p \) is achieved through modulating the pump beam in frequency and in intensity. That is, for a 5.0 mm long atomic lattice, the transmissivity peak reaches almost 100% at the third stop band edge and drops to near zero in less than 100 MHz. Correspondingly, the nonreciprocity parameter \( \Delta T_p \) approaches 96% for a nearly monochromatic pulse with the full bandwidth of 2.5 MHz (\( \mathcal{L}' = 120 \text{ m} \)). In Fig. 3 we plot the nonreciprocity parameter \( \Delta T_p \) as a function of both pump Rabi frequency \( \Omega_p \) and probe pulse detuning \( (\omega_p' - \omega_{31})/\gamma_{13} \). It is clear that the best nonreciprocity effect occurs when the pump Rabi frequency is set in the AT splitting regime with \( 10 \gamma_{13} \leq \Omega_p \leq 30 \gamma_{13} \).

Figure 4(a) further shows the dependence of the nonreciprocity parameter \( \Delta T_p \) on the pulse length \( \mathcal{L}' \). As we can see, \( \Delta T_p \) reduces from a value of about 92% to a value of about 38% for the sample velocity \( v = 15 \text{ m/s} \) when \( \mathcal{L}' \) is decreased from 120 m to 10 m. Figure 4(b) finally shows the dependence of the
We find that when $T_p$ approaches 100% for $L_0 = 120 \text{ m}$ its velocity dependence is nonlinear and tends to saturate already at $\approx 20 \text{ m/s}$. It is worth noting that the wide tuning of both magnitude and sign of $T_p$ achieved through the present reciprocity-breaking mechanism could be exploited to realize transmission modulation in hybrid optomechanical atom-membrane interfaces where the collective motion of ultracold atoms is strongly coupled to the vibration of a micromechanical membrane [32,33] (another similar means for enhancing an optomechanical coupling can be found in [41]). The unidirectional effect that we anticipate here for highly dispersive moving multilayers does not rely on any notion that is unique to atoms and hence the scheme may also be adapted to specific optomechanical metamaterial structures where the same effect could be observed or reversibly used for mechanical sensing [32,42,43]. On a more general ground, however, the reciprocity-breaking approach we suggest here is expected to stimulate further insights into the general transport theory of light in periodic structures [44], perhaps opening up novel possibilities in guiding nonreciprocity parameter $\Delta T_p$ on the sample velocity $v$. We find that when $\Delta T_p$ approaches 100% for $L' = 120 \text{ m}$ its velocity dependence is nonlinear and tends to saturate already at $v = 20 \text{ m/s}$.

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and redirecting the flow of electromagnetic radiation in such media.

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Note added in proof.—In proof-reading this article we have become aware of some very recent work that finds a related nonreciprocal effect in a spatially uniform distribution of atoms [46].

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[22] T is defined as the ratio of the transmitted pulse energy over the incident pulse energy.
[23] Conversely, when the transmissivity is a very steep function of frequency, which is the case of interest here, the pulse nonreciprocity parameter $\Delta T_p$ is to be calculated directly from (1).
[34] The atomic multilayer transverse radius $D = \sqrt{4Na/\pi Lnd}$ is determined by the in-wall atomic density $n = N/V$ and the number of atoms $N$ loaded into a sample of length $L$. For typical densities $n \sim 10^{12}$ cm$^{-3}$, filling factors $d/a = 0.1$ one obtains a transverse radius $D \sim 35 - 110 \mu m$ when $N \sim 10^8$ $10^9$ atoms are trapped in a sample of length $L \sim 1 cm$, which well justifies the 1D structure we examine here (see also [31]).
[44] The scheme—as devised here—is not suitable to break reciprocity in random photonic structures, where this may be achieved, e.g., through ultrafast nonlinear dephasing of light paths as shown in [45].