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Analysis of an all-optical SBS avalanche detector

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ABSTRACT

Seeding Brillouin scattering with a sufficiently efficient source of coherent phonons has the potential to produce energy-sensitive photon detectors. Based on this idea, we propose and analyze some possible designs for such a detector.

Keywords: Stimulated Brillouin scattering, optical fibers, detectors

1. INTRODUCTION

Some of the most revolutionary detectors have been based on preparing a medium in a state close to a phase transition where the non-equilibrium physics features dominate. Two relevant examples here are the Cloud Chamber[1] and its later counterpart, the Bubble Chamber.[2] Interestingly, both developments are based on optical readout in non-equilibrium media, in overcooled gases and superheated liquids, respectively. Here we consider the possibility of optical readout of non-equilibrium processes in solids.

Avalanche photodiodes operate on the principle of cascade amplification. A single photon is absorbed by the semiconductor material, producing an electron-hole pair; an applied electric field accelerates these charges, which collide with neutral atoms, creating more free charges through impact ionization. If at the initial stage, the electron and hole produced by photon absorption become trapped by impurities, the avalanche may not be triggered. Such losses result in a quantum efficiency less than one. In this paper, we consider an alternative approach, in which the first photon is initially converted into a very large number of phonons which may be subsequently amplified. As a large enough number of phonons is less likely to experience the same degree of loss as a single electron-hole pair, this process should have a higher quantum efficiency.

A suitable nonequilibrium process with very large gains is stimulated Brillouin scattering (SBS).[3,4] In SBS, pump and Stokes beams counterpropagate through a nonlinear medium. At their difference frequency, they drive acoustic vibrations in the medium through the process of electrostriction. When their frequency difference is close to the Brillouin frequency for the medium, their interaction creates a sustained acoustic wave: The pump beam, seeing a refractive index modulation that moves in the medium, scatters, causing the Stokes beam to be amplified and delayed. The envelope of the acoustic wave is driven by a term proportional to the product of the pump and Stokes fields.

Brillouin scattering is normally considered in two regimes: In the first regime, a strong pump amplifies a small initial Stokes pulse into a large one, and in the second, a strong pump interacts with thermal noise to produce Stokes light. We are interested in a third regime where rather than the acoustic field coming from thermal noise, a coherent acoustic field is deliberately injected into the medium. The effect of injecting coherent phonons into such a system is similar to injecting Stokes photons: Both result in an exponentially growing Stokes field which is paired with an exponentially increasing acoustic field: This combination we will refer to as a "Brillouin avalanche".

There is precedent for generating and using coherent phonons. In 1988, Rothenberg demonstrated the optical excitation of acoustic modes in a material by observing the motion of the surface of the material.[5] The field of picosecond ultrasonics has now developed this technique to the point where the optically excited and measured phonons can be used to acoustically probe a material sample.[6,7] Other groups have explored the creation of

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phonon avalanches: Amplification of phonons by stimulated emission between two Zeeman-split spin states has been demonstrated in low temperature ruby.[8,9]

Brillouin scattering in optical fibers has been used as the mechanism in a variety of detectors.[10,11] Most relevant is the work of Thévenaz, et al.[12] which showed how Brillouin scattering in a fiber continues to work at low temperatures (with a linewidth reduced by about a factor of 10 below a temperature of 10 K) and can be used as a cryogenic thermometer.

2. PRINCIPLE OF DETECTOR OPERATION

We begin with a single-mode fiber, pumped by monochromatic laser radiation from one end with frequency ω_p. Then in the stationary state, the pumping light intensity takes on the distribution $I_{\omega_p}(z) = I_{\omega_p}(0)\exp(-\alpha_p z)$, where z is the coordinate along the fiber, α_p is the absorption coefficient, and $I_{\omega_p}(0)$ is the intensity of the incident laser radiation. The fiber prepared in such a manner will amplify Stokes pulses at the frequency $\omega_s = \omega_p - \Omega_B$, where $\Omega_B \ll (\omega_s, \omega_p)$ is the Brillouin frequency of the fiber material. If a small Stokes pulse enters at $z = L$, while it propagates toward $z = 0$, it will interact with the pumping radiation and become amplified, creating an optical avalanche.

A similar (though quantitatively different) kind of avalanche can be observed without specially applied seeding pulses. Here the trigger is thermal fluctuations of the phonon field which can cause pump photons to be converted into seeding Stokes photons.[13] To use this mechanism as a detector of small amounts of energy (δE) absorbed from incident photons, one needs to reduce the noise level. This is a typical requirement for detectors operating on the principles of microcalorimetry. Noise reduction in these detectors is possible because their operation takes place in cryogenic environments. Deep cooling reduces the noise, and also makes the heat capacities of dedicated absorbers C_{abs} much smaller than at room temperature. Then the temperature excursion, $\delta T = \delta E / C_{abs}$ in these absorbers is measured with sensitive thermometers, which typically utilize front-end electronics (e.g., SQUID-array amplifiers). In one of our proposed detector designs, the fiber itself serves as absorber, thermometer and amplifier. In this case, the detecting end of the fiber can be deeply cooled because of the low thermal conductivity of glasses, and the temperature excursion can be appreciable since glasses generally have very low specific heat, so δT will be relatively big.

The other (more efficient, but also, more challenging) option is to convert the energy of the photons to be detected into coherent phonons which are then injected into an optical fiber where they trigger an optical avalanche. Unlike avalanche photodiodes, which typically operate in saturated mode, an avalanche based on Brillouin scattering can be held to within the linear amplification regime and therefore has the potential to resolve photon energies. In this article, we will consider the fundamentals of such detectors and discuss possible approaches to realizing them.

3. BASIC FORMALISM

Following the general description of Brillouin scattering in a single-mode fiber[3,4,14] and neglecting optical losses and higher-order derivatives, one can write the equations for the slow varying amplitudes of the pumping field (E_p), Stokes field (E_s) and acoustic field (ρ) in the form:

$$\frac{\partial E_s}{\partial z} + \frac{n}{c}\frac{\partial E_s}{\partial t} = \frac{i\gamma\omega_s}{4\rho_0 nc}\rho^* E_p \tag{1}$$

$$\frac{\partial E_p}{\partial z} - \frac{n}{c}\frac{\partial E_p}{\partial t} = \frac{i\gamma\omega_p}{4\rho_0 nc}\rho E_s \tag{2}$$

where n is the refractive index of the material at the pump frequency, $\gamma = \rho_0 (d\varepsilon/d\rho)$ is the electrostrictive constant, and ρ_0 is the average density of the fiber material.

The most fundamental and unavoidable noise source in SBS is caused by thermal fluctuations of the fiber material. (We will neglect the noise contributions of Rayleigh scattering, as the spontaneous Stokes noise will dominate when the pump power is above the threshold for SBS.[15] We also note that Rayleigh scattering is elastic in nature and might potentially be filtered out.) This noise channel is usually modelled by including a Langevin noise term f in the differential equation for the fiber material density variations, ρ:[16]

$$\frac{\partial \rho}{\partial t} + \frac{\Gamma}{2}\rho = i\Lambda E_p E_s^* + f. \tag{3}$$

where Γ is the acoustic damping rate and where Λ is determined by the electrostrictive strength. The function $f = f(z,t)$ corresponds to a Gaussian random process which has zero mean and is δ-correlated in space and time:

$$\langle f(z,t) f^*(z',t') \rangle = Q\delta(z - z')\delta(t - t'), \tag{4}$$

where the amplitude Q is obtained from the fluctuation-dissipation theorem to be

$$Q = 2k_B T \rho_0 \Gamma / \left(v_A^2 A_{eff}\right). \tag{5}$$

Here A_{eff} is the effective mode area and v_A is the acoustic velocity at the Brillouin frequency in the medium, k_B is Boltzmann's constant, and T is the temperature of the medium.

4. SIGNAL VS. NOISE FOR THE ACOUSTIC FIELD

The noise properties of such a detector depend sensitively on the nature of the mechanism for converting energy into the phonon modes for Brillouin amplification. However, there are two limiting cases: In the thermal case, the energy δE of the detected photon is transfered into thermal energy, briefly raising the noise floor and initiating a temporary increase in spontaneous Stokes generation (effectively creating a Stokes pulse). In the case of resonant energy transfer, the energy δE is transfered into the required Brillouin phonon mode thus triggering the Stokes pulse. Let us consider them in detail.

4.1 Thermal limit

The temperature deviation in this case is

$$\delta T = \frac{\delta E}{C_{abs}} \tag{6}$$

where $C_{abs} = c_V V_{abs}$ is the heat capacity of the absorber, c_V is the specific heat, and V_{abs} is the heat absorber volume, which in the optimal case coincides with that of the fiber core. Note that the heat capacity strongly depends on temperature, becoming very small at low temperatures, thus enhancing the signal output. Our estimates suggest that this approach would only begin to be plausible at temperatures at or below 0.1 K. We note that this is a typical temperature for cryogenic microbolometers, so a detector based on conversion of incident energy into thermal phonons may be worth considering. However, for the remainder of this article, we will focus on the possibility of the detection of photons via conversion to coherent phonons.

4.2 Resonant limit

In this case, we assume that the total deposited energy δE can be converted with efficiency η $(0 \le \eta \le 1)$ into the phonon mode Ω_k in the core of the optical fiber, creating the nonequilibrium phonon population with occupation number $N_k = \eta\delta E/\hbar\Omega_k$. To convert this value into the density perturbation $\delta\rho(z,t)$, we use the quantum-mechanical relation:[17]

$$\delta\widehat{\rho}(z,t) = \sum_k i \left(\frac{\hbar k \rho_0}{2V v_A}\right)^{1/2} \left(\widehat{c}_k e^{i(kz - v_A k t)} - \widehat{c}_k^{+} e^{-i(kz - v_A k t)}\right) \tag{7}$$

where the creation $(\widehat{c}_k^{+})$ and annihilation $(\widehat{c}_k)$ operators are related with the phonon occupation number by the expression:

$$\langle N_{k-1}|\widehat{c}_k|N_k\rangle = \langle N_k|\widehat{c}_k^{+}|N_{k-1}\rangle = \sqrt{N_k}. \tag{8}$$

If $N_k \gg 1$ (*i.e.*, $\hbar\Omega_k \ll \delta E$, which is typically the case), then all these operators may be replaced by c-numbers, and if only one mode k is essential in Eq. 8, then (taking mode volume $V = V_{abs}$),

$$|\delta\rho| \approx \left(\frac{\hbar k \rho_0}{2 V_{abs} v_A} N_k\right)^{1/2}. \tag{9}$$

We estimate the signal-to-noise ratio, R from Eq. 3. We assume that the noise term dominates over the coherent driving term (at least at $t = 0$) when the energy absorption takes place. Then, at a given position, z, Eq. 3 converts into a Langevin equation of the Ornstein-Uhlenbeck theory:[18]

$$m\frac{\partial u(t)}{\partial t} = -\xi u(t) + \lambda(t) \tag{10}$$

for a particle with a mass m and friction in a phase space $(x, u; u \equiv dx/dt)$ with the designated white noise driving force $\langle\lambda(t)\lambda(t')\rangle = 2\xi k_B T \delta(t - t')$. Subject to a Maxwell-Boltzmann distribution of velocities, the stationary solution of Eq. 10 is

$$\langle u^2 \rangle = k_B T/m. \tag{11}$$

In order to use the analogous reasoning on Eq. 3, we need to account for the z-dependence of the Langevin noise term $f(z,t)$, (Eq. 4). Normally, one would perform a double integral over the z and z' part of f, each of which contains the square root of the Dirac delta function. The double integral has a $1/L_{abs}$ coefficient, so that the total double integral is normalized. However, in our case we are evaluating $\langle\rho^2\rangle$ at a particular value of z, so we just integrate over z' to obtain:

$$\frac{1}{L_{abs}} \int_0^{L_{abs}} \delta(z - z')\, dz' = \frac{1}{L_{abs}} \qquad (\text{for } 0 \leq z \leq L_{abs}) \tag{12}$$

Now, comparing Eqs. 10 and 3, we find that the correct form of $\langle\rho^2\rangle$ is:

$$\langle\rho^2\rangle = \frac{Q}{\Gamma L_{abs}}. \tag{13}$$

This should be compared with the squared deviation of ρ, created by the energy deposition δE (Eq. 9). The resulting signal-to-noise ratio is

$$R = \frac{\hbar k \rho_0 N_k}{2 V_{abs} v_A}\frac{\Gamma L_{abs}}{Q} = \eta\frac{\delta E}{4 k_B T} \tag{14}$$

Setting the signal-to-noise ratio equal to 1, we then get the smallest resolvable energy to be: $\delta E = 4k_B T/\eta$, where $L_{abs} = V_{abs}/A_{eff}$ is the effective length of fiber in which the energy was deposited. In the ideal case of converting all incident energy into phonons ($\eta = 1$), energies on the scale of a few $k_B T$ could locally create a modulation in the acoustic field, resolvable from thermal noise.

This is an upper bound, neglecting losses and the addition of noise associated with transforming the energy δE into phonons. If it were possible to read this out, it would create a significant technological advance in energy resolution.

5. SIGNAL VS. NOISE FOR THE STOKES FIELD

In order to calculate the signal-to-noise ratio for the Stokes output, we use the analytical results on Brillouin scattering from Boyd, et al.[14] When seeded by thermal noise, the average Stokes field squared at the end of a fiber is

$$< |E_s(0)|^2 >= \left(\frac{\gamma\omega_s}{4\rho_0 n c}\right)^2 |E_p|^2 \frac{QL}{\Gamma} e^{G/2} \left[I_0(G/2) - I_1(G/2)\right] \tag{15}$$

where $I_0(x)$ and $I_1(x)$ are modified Bessel functions of the first kind.

For large values of G ($G > 21$), the amount of pump power being reflected back (as spikes in the thermally triggered Stokes) is large, and this limits signal-to-noise ratio and dynamic range (see, for instance, Fig. 4 of Ref. 14). To avoid this condition, we choose to limit our consideration to values of $G < 21$.

From Ref. 14 we have an expression for the amount of Stokes intensity generated by thermal noise. Suppose we excite N_k coherent phonons at the end of the fiber over a length ΔL. (We will refer to this as a step-function distribution.) For simplicity, we assume that the field from the coherent phonons is large enough that the thermal contributions may be ignored. Integrating the Stokes equation $\partial E_s/\partial z = -i\Psi\rho_L^* E_p$ (where $\Psi = \gamma\omega_s/4\rho_0 nc$) and neglecting back-action on the pump field, we get $E_s(L) = i\Psi\rho_L^* E_p \Delta L$ and so $E_s(0) = i\Psi\rho_L^* E_p \Delta L \exp(G/2)$.

The signal-to-noise ratio is then

$$R = \frac{|\rho_L|^2 (\Delta L)^2 e^{G/2}\Gamma}{QL\left(I_0(G/2) - I_1(G/2)\right)} \tag{16}$$

In the limit of large G, $I_0(G/2) - I_1(G/2)$ becomes $\exp(G/2)/\left(\sqrt{\pi}G^{3/2}\right)$, and the signal-to-noise ratio reduces to

$$R = \frac{\sqrt{\pi}|\rho_L|^2 (\Delta L)^2 \Gamma G^{3/2}}{QL} \tag{17}$$

This should certainly be valid for large values of R. Suppose it still holds down to $R = 1$. (It seems reasonable as a lower bound on the required number of phonons.) This then tells us the necessary $|\rho_L|^2 \Delta L$ product to obtain a given signal-to-noise ratio:

$$|\rho_L|^2 = \frac{RQL\left(I_0(G/2) - I_1(G/2)\right)}{(\Delta L)^2 e^{G/2}\Gamma}\Big|_{\text{large } G} \approx \frac{RQL}{\sqrt{\pi}(\Delta L)^2 \Gamma G^{3/2}} \tag{18}$$

Then we can use Eq. 9 to convert this into an expression for the corresponding number of phonons.

$$N_k = \frac{2R\left(I_0(G/2) - I_1(G/2)\right)}{e^{G/2}} \frac{A_{\mathit{eff}} L v_A}{\Delta L \Gamma} \frac{Q}{\hbar k \rho_0} \approx \frac{2R}{\sqrt{\pi}G^{3/2}} \frac{A_{\mathit{eff}} L v_A}{\Delta L \Gamma} \frac{Q}{\hbar k \rho_0} \tag{19}$$

Substituting in the expression for Q from Ref. 14,

$$Q = \frac{2\rho_0 \Gamma \hbar \Omega (1 + \bar{n})}{v_A^2 A_{\mathit{eff}}}, \tag{20}$$

we get

$$N_k = \frac{8R\left(I_0(G/2) - I_1(G/2)\right)}{e^{G/2}} \frac{L}{\Delta L}(1 + \bar{n}) \approx \frac{8R}{\sqrt{\pi}G^{3/2}} \frac{L}{\Delta L}(1 + \bar{n}) \tag{21}$$

By similar manipulations, we can express the signal-to-noise ratio in terms of the ratio of coherent phonons to thermal phonons:

$$R \approx \frac{N_k}{(1 + \bar{n})} \frac{\Delta L}{L} \frac{G^{3/2}\sqrt{\pi}}{8} \tag{22}$$

or transforming back via $\hbar\Omega(1 + \bar{n}) = k_B T$ and $\hbar\Omega N_k = \delta E \eta$,

$$R \approx \eta \frac{\delta E}{k_B T} \frac{\Delta L}{L} \frac{G^{3/2}\sqrt{\pi}}{8} \tag{23}$$

With high gain chalcogenide fiber, operated with a G of 20, this signal-to-noise ratio could be improved over that of Eq. 14. For an incident detected energy of $\delta E = 1.3$ eV (corresponding to a photon with wavelength 1 μm), the signal-to-noise ratio's dependence on temperature is as shown in Fig. 1. We assume that the ratio of ΔL to L is 0.1.

The efficiency, η will be determined by three factors: 1) input coupling (which can be near unity as for optical cavity modes), 2) photon-phonon conversion efficiency, 3) coupling of phonons to the Brillouin amplification stage.

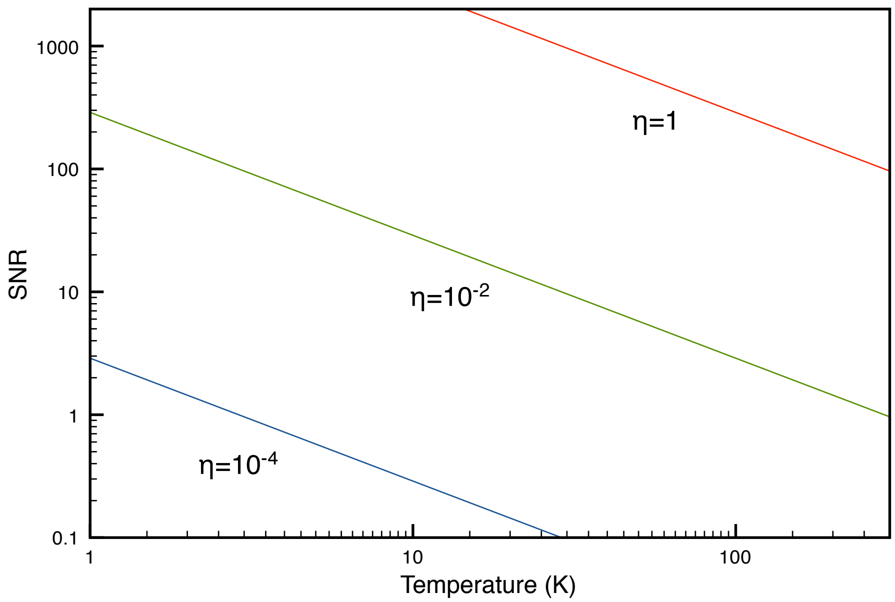

Figure 1. Signal-to-noise ratio, as specified in Eq. 23, versus temperature, evaluated at an incident energy of $\delta E = 1.3$ eV, $\Delta L/L = 0.1$, and $G = 20$, and for several values of η.

5.1 Signal vs. noise in the phonon production channel

A more precise calculation would consider the diminished phase coherence associated with converting a single photon into multiple phonons, as required by unitary evolution. An ideal, linear, phase-insensitive amplifier with gain N has a noise figure of $N/(2N-1)$,[20] so a rough estimate for the best-case scenario is that for large N, the signal-to-noise ratio will be reduced by 3 dB in the process of converting a single photon into N phonons. The N phonons will therefore not be perfectly coherent. A reduction in phase coherence of the phonons will decrease the coherence of the generated Stokes photons, reducing overall detector efficiency.

We also note that the transformation of the energy of the initially absorbed photon into the energy of secondary quasiparticles (in our case, acoustic field quanta) is subject to statistical (typically, Poisson) uncertainty, which implies inherently unavoidable fluctuational contribution (shot noise), and hence sets additional restrictions to the energy resolution of the detector. Keeping in mind that in the ideal case, the absorbed energy, E creates $N_{ph} = \eta E/\hbar\omega_B$ phonons, and that the uncertainty in case of Poisson statistics is $\Delta N_{ph} = \sqrt{N_{ph}}$, we have $\Delta E/E = \Delta N_{ph}/N = \sqrt{\hbar\omega_B/(\eta E)}$. This result for the energy resolution is commonly written in the canonical form (see, e.g., Ref. 21):

$$\Delta E_{FWHM} = 2.355\sqrt{FE\hbar\omega_B/\eta} \tag{24}$$

where F (the Fano factor) reflects the fact that in reality the production of the phonons is not fully uncorrelated (shot noise suppression in semiconductors typically yields $F \sim 0.1-0.2$) and where the index "FWHM" indicates that a standard definition of energy resolution is being used, instead of assuming that the signal-to-noise ratio is 1; this definition introduces the factor $2\sqrt{2\ln 2} = 2.355$ as it assumes a Gaussian distribution. It is important for our consideration that for the Brillouin phonons, the energy price per secondary quasiparticle, $\hbar\omega_B$, is at least five orders of magnitude smaller than the price per secondary quasiparticle in semiconductors (the energy gap, E_g) and two orders of magnitude smaller than the energy gap in superconductors (typically denoted as 2Δ). This provides us grounds to conclude, in view of Eq. 24, that the statistical uncertainty channel taken alone promises an order of magnitude better resolution for Brillouin avalanche detectors than superconducting tunnel junction detectors and two orders of magnitude better resolution than semiconductor-based detectors.

When combining various noise sources, different noise types at any given stage typically add in quadrature. For small phonon numbers, shot noise will dominate. Otherwise, thermal noise is expected to be the limiting factor, though depending on the nature of the transducer, photon-to-phonon conversion noise could prove stronger.

Our signal-to-noise rate calculations only provide a crude first estimate of the noise properties of Brillouin amplification as they use an averaged noise power. Ultimately, a detector is limited by the large but rare spikes in the noise which are better characterized by a dark count rate.

5.2 An exact, closed form solution

Instead of spatially separating regions of Brillouin gain from regions where coherent phonon seeding takes place, we can use the Raymer integral (Eq. 19 from Ref. 19) for a unified treatment of noise, acoustic seeding, and photonic seeding. This integral is an exact solution of the SBS equations (when higher-order derivatives have been neglected and the pump can be assumed to be undepleted). In our formalism, this equation for the Stokes field becomes:

$$\begin{aligned} E_s(z,\tau) = E_s(L,\tau) + \sqrt{\frac{\Gamma G(L-z)}{4L}} \int_0^\tau d\tau' e^{-\Gamma(\tau-\tau')/2} \frac{E_s(L,\tau')}{\sqrt{\tau-\tau'}} I_1\left(\sqrt{\frac{G\Gamma(\tau-\tau')(L-z)}{L}}\right) \\ + \frac{i\gamma\omega_s}{4\rho_0 nc} E_L \int_0^\tau d\tau' \int_z^L dz' e^{-\Gamma(\tau-\tau')/2} f^*(z',\tau') I_0\left(\sqrt{\frac{G\Gamma(\tau-\tau')(z'-z)}{L}}\right) \\ + \frac{i\gamma\omega_s}{4\rho_0 nc} E_L e^{-\Gamma\tau/2} \int_z^L dz' \frac{s(z',0)}{2} I_0\left(\sqrt{\frac{G\Gamma\tau(z'-z)}{L}}\right) \end{aligned} \quad (25)$$

Here, $\tau = t - zn/c$ is the local time. The initial acoustic field is $s(z,0)$. Note also that there is no cross-coupling between the acoustic seed and the Stokes seed in this model; they evolve independently and add linearly to the Stokes field. We wish to optimize the conversion of phonons into Stokes photons. As the constrained quantity is the acoustic field squared, integrated over space, and as the expression for the number of Stokes photons requires squaring the Stokes field and integrating over time, the optimum initial acoustic distribution is not immediately apparent. Our numerical calculations suggest that it is dependent on SBS gain and is not as simple as placing all the phonons at one end of the fiber or distributing them evenly over the fiber.

The best step-function distribution deposits phonons in the last ∼25 centimeters of the fiber. This is not unexpected, since the phonons have a finite lifetime (on the order of 10 ns), and the distance that an optical pulse can travel during that time is on the order of 30 centimeters. The filtering provided by the resulting 25 MHz Brillouin linewidth causes a Stokes pulse to experience decreasing gain as its duration is reduced. As a consequence, bursts of noise that are short compared to the acoustic lifetime will contribute only weakly to Brillouin scattering, yielding reduced noise for detector applications.

If we take the same number of phonons used in the step-function distribution and rearrange them into a distribution that decays exponentially from $z = L$ toward $z = 0$, we will obtain a 50% enhancement in Stokes output. Varying the phases of the fields may also help. The work of Thévenaz[22] suggests that there may be benefits to flipping the phase of the pump beam briefly. What the shape of the optimal distribution is, is still an open question.

6. POSSIBLE PRACTICAL REALIZATIONS

One of the keys to designing highly sensitive detectors based on our Brillouin avalanche scheme is converting the incident photon into sound. We will consider here some possible approaches.

1a) One can conceive of lithographically-manufactured Brillouin transducers which transform the acoustic bursts from the energy deposition of a single-photon into the geometro-acoustic modes in crystalline plates which will be designed to be in exact resonance with the known Brillouin mode of the waveguide. The principle of such a transducer is shown in a simple form in Fig. 2.

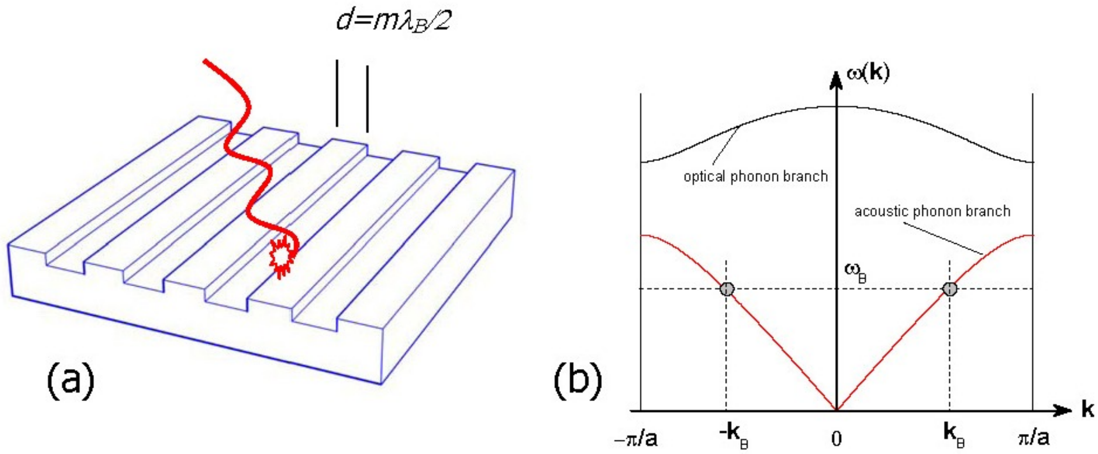

Figure 2. High-energy photon absorption causes an impulsive blast which has a wide Fourier spectrum that will spread above $\omega_B/2\pi \sim$ 10GHz. If the absorber has geometric structure, for example, periodic structure (panel a) facilitating resonance with Brillouin modes at $\omega_B = 2\pi v_A/\lambda_B$, coherent vibrations at this mode will build up in the whole structure. If this material is acoustically coupled with the read-out cell (see below), acoustic phonons with wave vectors k_B and $-k_B$ (panel b) will be induced in the cell.

1b) Optical cells that contain planar fiber structure made of highly nonlinear material (such as chalcogenide glass) could also be manufactured by lithographic methods. These cells can be integrated into fiber-optic circuits for observation of stimulated Brillouin scattering in the cells.

To join these two components in the final design, the transducer and the cell could be manufactured on the opposite surfaces of a single crystalline plate, thus comprising a simple design for the Brillouin avalanche detectors.

2a) Another approach is to confine the photon to be detected in a high-Q cavity. Recent advances in microresonator structures have achieved significant improvements in optomechanical coupling strength. Noise reduction in such structures is possible, both by conventional cooling and more targeted cooling. Several groups have recently demonstrated sideband cooling (combined with cryogenic cooling of the bath to temperatures of 1-5 K) of a microresonator so that the mode of interest only has 30 to 60 acoustic quanta in it.[23–25]

What is the maximum rate that we could expect to extract phonons from light trapped in a microresonator? Eichenfeld, et al.[26] argue that light in a Fabry-Perot cavity will transfer 2 quanta of momentum to the 2 mirrors for each round trip. For light with a wavelength of 1 micron, a photon will make a round trip in a cavity of length λ in ∼6 fs. Therefore the maximum phonon generation rate is 1 phonon every 3 fs or 3×10^5/ns. Most microresonators appear to have an optomechanical coupling constant that also scales according to the path length of the light (possibly a factor of 2π better).

2b) The zipper cavities of the Painter group[26] are unique designs that can exceed this path-length-dependent optomechanical coupling strength. They serve as simultaneous cavities for photons and phonons and have, in the earliest design, shown sensitivity to the cavity motion induced by 660 photons.[26] A more advanced design has allowed smaller masses and scales to be achieved and has now made possible transfer of photon momentum to the cavity over a length close to the wavelength of the light.[27] Further, they can operate in air and at room temperature, making this a strong candidate for practical, highly efficient conversion between light and sound.

We note that, particularly with this approach, there is a connection between our proposal and a recently proposed scheme for counting microwave photons by using nonlinear optics to convert them to the optical regime, using a resonator for enhancing both the microwave and optical signals.[28]

3) A drawback to using an optical cavity (as with the zipper cavity) is that the optical resonances will only allow the detector to work with high efficiency for a small set of frequencies. A more speculative alternative

to an optical cavity would be the recently proposed "optical black hole".[29] The proposal prescribes using a continuously variable dielectric permittivity to construct a potential analogous to a gravity well. Essentially it calls for a smooth reduction in refractive index from a larger value outside to a smaller value in some inner core region. This allows light to be effectively pulled in and then trapped by the index gradation over a wide range of frequencies. This technique has since been demonstrated at microwave frequencies by constructing a decreasing refractive index (to values below that of air) from metamaterials,[30] and it is expected that this approach will soon be demonstrated at optical wavelengths.

While this system may be considered more analogous to an ideal black body (as the trapped light is eventually converted to heat, which is then radiated away), the broadband trapping should enable the capture of nearly all incident light. Combining such high-efficiency photon capture with a mechanism for extracting energy from the photon in the form of phonons would allow efficient optoacoustic conversion. Embedding such an acoustic transducer in the optical fiber, by using appropriately designed microstructure in the fiber, could provide maximum conversion efficiency.

The final missing factor in the optoacoustic conversion efficiency, η is the coupling of phonons to the acoustic mode of the optical fiber which will depend on the transducer scheme and geometry. Determining the best approach and the resulting overall efficiency represents the biggest remaining challenge to implementation of our scheme.

7. DISCUSSION AND CONCLUSIONS

In APDs, the photon initially creates only one electron-hole pair. In avalanche Brillouin detectors (ABDs), the incident photon could create multiple initial phonons, up to $\sim 10^5$ for visible light (approaching, in the first amplification stage, the total amplification of a typical APD). Typically, when the amplification reaches the value of a million in avalanche photodetectors (APDs), they become effective as single-photon avalanche photodetectors (SAPDs). Brillouin amplification of a billion is readily attainable.

Single-photon counting detectors, such as superconducting single-photon detectors[31] are now highly developed, with good dark count rates and high quantum efficiency. However, as with other highly sensitive avalanche detectors, they are energy blind. The reason why our proposed detector is different is that it is biased in the unsaturated mode, so that the optical output is still proportional to the energy input, making it at least conceivable that this scheme might provide a photon-counting, energy-discriminating detector. Such detectors would find application in quantum computing for preparing Fock states. In principle, the dynamic range of this detector could be easily varied by changing the intensity of the pumping beam, adapting the sensitivity to detect smaller or larger energies.

We note that the generality of this scheme allows it to be adapted to detection of particles rather than photons, should a sufficiently efficient particle-to-phonon transducer be found. Microcantilever technology may provide a useful way to convert incident kinetic energy to coherent phonons.

Although we attempted to adequately model noise in this paper, it is well-known that SBS is an inherently noisy gain process. As there are still unaddressed mechanisms and parameters that need to be measured, we do not have the expectation that our scheme can detect single optical photons but believe that it may be useful for detecting low-light levels (small numbers of optical photons). Another possibility is to convert high frequency photons, such as X-rays, to a sufficiently large number of phonons that the thermal noise can be overcome. If a highly efficient optoacoustic transducer can be made, in the right regime, Brillouin avalanche detectors may become useful devices.

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