Editors' Suggestion

Stimulated Emission or Absorption of Gravitons by Light

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(Received 3 March 2025; revised 25 April 2025; accepted 26 August 2025)

We study the exchange of energy between gravitational and electromagnetic waves in an extended Mach-Zehnder or Sagnac type geometry that is analogous to an "optical Weber bar." In the presence of a gravitational wave (such as the ones measured by the Laser Interferometer Gravitational Wave Observatory), we find that it should be possible to observe (via interference or beating effects after a delay line) signatures of stimulated emission or absorption of gravitons with present-day technology. Apart from marking the transition from passively observing to actively manipulating such a natural phenomenon, this could also be used as a complementary detection scheme. Nonclassical photon states may improve the sensitivity and might even allow us to test certain quantum aspects of the gravitational field.

DOI: 10.1103/xd97-c6d7

Introduction—In the history of electrodynamics, an important step was Franklin's pioneering (though extremely dangerous) kite experiment, where static electricity was collected from the air by flying a kite into or close to thunder clouds. On the one hand, this experiment showed that lightning and electricity, as known from laboratory experiments (e.g., with Leiden jars or by rubbing amber), are basically of the same nature and has thereby made a significant contribution to unifying these phenomena—eventually leading to our modern understanding of electrodynamics and the standard model. On the other hand, the kite experiment marked the transition from passively observing a natural phenomenon such as lightning to actively manipulating it—and thereby paved the way for many modern technological developments, from lightning rods to power plants.

Owing to the weakness of the gravitational interaction (in laboratory scale experiments) as determined by Newton's constant, $G_{\rm N} \approx 6.7 \times 10^{-11}\ {\rm m^3\,kg^{-1}\,s^{-2}}$, we are now in a somewhat similar situation regarding gravitational waves. They were predicted by Einstein around a century ago [1,2]. However, it took more than half a century before indirect evidence for them was observed in the Hulse-Taylor binary pulsar [3,4], whose energy loss over time due to the emission of gravitational waves agrees very well with the predictions from general relativity. Recently, the direct detection of gravitational waves on Earth was achieved at the Laser Interferometer Gravitational Wave Observatory (LIGO) [5,6]. Both accomplishments mark important breakthroughs; they were awarded with the Nobel Prizes in Physics in 1993 and 2017, respectively.

In the following, we shall make the assumption (which is quite natural but has not yet been verified experimentally) that gravitational waves are, at least in the weak-field regime, analogous to electromagnetic waves in the sense that their energy is quantized in terms of excitation quanta $\hbar\omega$ (i.e., gravitons), where ω is the frequency of the gravitational wave. Possible consequences of departure from this assumption will be discussed below. Then, to facilitate the transition from passively observing a natural phenomenon such as gravitational waves to actively manipulating it, let us ask the following question: Can one design an experiment where at least one graviton with energy $\hbar\omega$ is emitted (or absorbed) in a verifiable manner? Or, to phrase it in another way, Can we recreate a Hulse-Taylor-like scenario in the laboratory? As a first approach to this question, let us take the well-known quadrupole formula describing the power P emitted by gravitational radiation (in analogy to the dipole formula in electromagnetism)

$$P = \frac{G_{\rm N}}{45c^5}\sum_{ij}\left(\dddot{Q}_{ij}\right)^2, \tag{1}$$

where Q_{ij} are the quadrupole moments of the dynamical mass distribution [7]. In terms of its characteristic length L and mass m, they scale as $Q_{ij} = \mathcal{O}(mL^2)$. Thus, the total emitted power is $P = \mathcal{O}(\omega^6 m^2 L^4 G_{\rm N}/c^5)$, where ω is the oscillation frequency (i.e., the frequency of the emitted gravitational waves). Together with Newton's constant $G_{\rm N}$, the speed of light $c \approx 3 \times 10^8\ {\rm m/s}$ suppresses the prefactor

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in the quadrupole formula (1) by more than 50 orders of magnitude when expressed in terms of SI units; see also [7].

Thus, even after comparison to the small Planck constant $\hbar \approx 10^{-34}$ J s, we find that it will be extremely difficult to emit one gravitational excitation quantum (graviton) with energy $\hbar\omega$ using everyday values of m, L, and ω in the kilogram, meter, and second (hertz) regime [8]. To cast the result into a dimensionless form, let us introduce the number of gravitons N emitted per oscillation period $N = \mathcal{O}(P/[\hbar\omega^2])$, the characteristic velocity scale $v = \omega L$, and the Planck mass $m_{\mathrm{P}} = \sqrt{\hbar c/G_{\mathrm{N}}} \approx 22$ μg. Then we find that N scales as m^2/m_{P}^2 multiplied by v^4/c^4, showing a strong suppression for slow velocities, i.e., in the nonrelativistic regime.

These considerations suggest using light [9]. Since stationary plane-wave cw laser beams do not emit gravitational waves [10], we consider laser pulses in suitable geometries; see also [11]. As an extreme example, let us take the megajoule pulses at the National Ignition Facility [12]. Still, they correspond to a mass of order 10^{-11} kg, i.e., well below the Planck mass. As a result, a rough order-of-magnitude estimate obtained by combining this mass of order 10^{-11} kg (squared) with Newton's constant G_{N} in comparison to $\hbar c \approx 3 \times 10^{-26}$ J m already shows that it is still very difficult to emit a single graviton in this way.

Hence, we follow a different route here and consider the stimulated emission of gravitons instead of their creation by quadrupole formula (1); see also [13–15] (though in a different context). To this end, we consider light pulses propagating within a preexisting gravitational wave (such as the ones measured by LIGO, for example) and determine the transfer of energy between the gravitational and the electromagnetic field; see also [16]. Finding an energy transfer of $\hbar\omega$ or more is then interpreted as a smoking gun for the emission or absorption of gravitons by light. Note that this scheme displays some similarities to resonant mass antennas such as Weber bars [17–19], but since we are using a highly excited state (a light pulse), we may not only absorb but also emit gravitational radiation.

Gravitational waves—For simplicity, we consider linearly polarized gravitational waves propagating in the z direction, but our results can be generalized to other wave forms in a straightforward way. In a suitable coordinate system, the metric reads ($\hbar = c = \varepsilon_0 = \mu_0 = 1$)

$$ds^2 = dt^2 - [1+h]dx^2 - [1-h]dy^2 - dz^2, \qquad (2)$$

where $h(t,z) = h(t-z)$ is the amplitude of the gravitational wave. Since this quantity is extremely small, for example, $h = \mathcal{O}(10^{-22})$, we neglect second and higher orders in the following. Thus, the metric determinant simplifies to $\sqrt{-g} = 1 + \mathcal{O}(h^2)$. Furthermore, in view of the long wavelength of the gravitational waves and the fact that we consider light pulses propagating in the x-y plane, we may neglect the spatial dependence $h(t,z) \approx h(t)$.

In these coordinates [Eq. (2)], the Christoffel symbols corresponding to Newton's gravitational acceleration vanish ($\Gamma^i_{00} = 0$), and thus massive objects such as the mirrors used to reflect the light pulses stay at rest. However, since the x and y coordinates are rescaled differently by the gravitational wave, it could affect the angle under which the light pulses are reflected. In principle, this angular deflection of order $\mathcal{O}(h)$ could also be used to detect gravitational waves. However, since laser beams and pulses with a well-defined propagation direction must have a sufficiently large width (of many wavelengths), the impact of this tiny deflection angle can be neglected here.

Electromagnetic waves—Now let us consider light pulses propagating in the background [Eq. (2)]. To maximize the effect, we consider light polarized in the z direction $\boldsymbol{A}(t,\mathbf{r}) = A_z(t,x,y)e_z$, but again our analysis can easily be generalized. Note that the form $\boldsymbol{A}(t,\mathbf{r}) = A_z(t,x,y)e_z$ automatically satisfies the generally relativistic Lorenz gauge condition $\nabla_\mu A^\mu = 0$. The contraction $F_{\mu\nu}F^{\mu\nu}$ of the field strength tensor $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ gives the Lagrangian density

$$\mathcal{L} = \frac{1}{2}\left\{(\partial_t A_z)^2 - [1-h](\partial_x A_z)^2 - [1+h](\partial_y A_z)^2\right\}. \qquad (3)$$

Field quantization yields the interaction Hamiltonian

$$\hat{H}_{\mathrm{int}} = h\int d^3r\left[(\partial_y\hat{A}_z)^2 - (\partial_x\hat{A}_z)^2\right], \qquad (4)$$

which is determined by the magnetic fields $\hat{B}_x^2 - \hat{B}_y^2$ and describes the coupling between the electromagnetic field and the gravitational wave.

Now we can study the energy transferred between these two. To this end, we employ the Heisenberg picture with $d\hat{H}/dt = (\partial\hat{H}/\partial t)_{\mathrm{expl}}$ where the explicit time dependence stems from the gravitational wave, i.e., $h(t)$. Taking expectation values then yields the energy transfer

$$\frac{d\langle\hat{H}\rangle}{dt} = \dot{h}\int d^3r\left\langle(\partial_y\hat{A}_z)^2 - (\partial_x\hat{A}_z)^2\right\rangle, \qquad (5)$$

where the divergent vacuum contributions $\langle 0|(\partial_y\hat{A}_z)^2|0\rangle$ and $\langle 0|(\partial_x\hat{A}_z)^2|0\rangle$ cancel each other such that we can use renormalized (e.g., normal ordered) values. The integral on the right-hand side of Eq. (5) is the difference between the total energies of the light pulse in the magnetic field components in the x and y directions. Thus, we find a rigorous bound $|\dot{E}| \le |\dot{h}|E$ for the energy transfer $\dot{E} = d\langle\hat{H}\rangle/dt$ in terms of the total energy E of the laser pulse [20]. In practice, however, it is very difficult to saturate this bound since the light energy oscillates rapidly

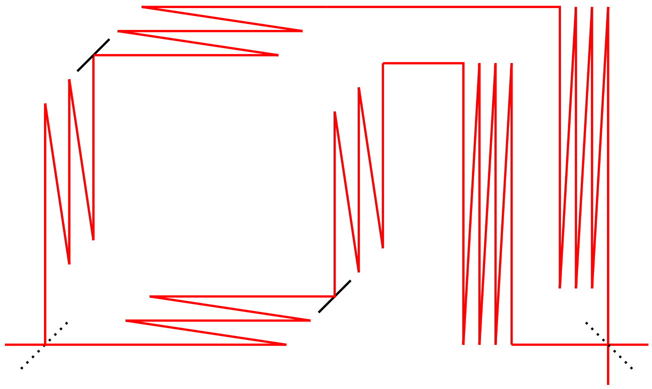

FIG. 1. Sketch (not to scale) of a possible geometry. The initial laser pulse is split up by a half silvered mirror (dotted black line on bottom left) into two pulses (red lines), first propagating in the x and y directions, respectively. After half a period (ideally at $\dot{h} = 0$), these pulses are reflected by the 45° mirrors (solid black lines) in order to propagate in the respective other directions. After traversing this Mach-Zehnder or Sagnac type geometry on the left-hand side and thereby gaining or losing energy, the light pulses are sent through further optical paths of equal length (or basically the same path) on the right-hand side in order to accumulate a large enough phase difference. Finally, they are brought to interference at another half silvered mirror (dotted black line on bottom right). The optical paths are elongated by retroreflections. The mirrors for doing that and for guiding the pulses are omitted for simplicity.

between the electric component $\langle(\partial_t\hat{A}_z)^2\rangle_{\rm ren}$ and the magnetic components $\langle(\partial_x\hat{A}_z)^2\rangle_{\rm ren}$ or $\langle(\partial_y\hat{A}_z)^2\rangle_{\rm ren}$. Thus, we have $|\dot{E}| \leq |\dot{h}|E/2$ on average.

To maximize energy transfer, one could imagine the following scenario (see also Fig. 1). As long as $\dot{h} > 0$, we have a light pulse propagating in the x direction, and then—after reflection with a mirror—it propagates in the y direction as long as $\dot{h} < 0$, and so on. In this case, we have $\dot{E} < 0$ and thus the emission of gravitons (in view of energy conservation). The opposite case (the x direction for $\dot{h} < 0$ and the y direction for $\dot{h} > 0$) yields $\dot{E} > 0$ and thus the absorption of gravitons.

Experimental parameters—Let us study the experimental feasibility of the above scheme by inserting typical example values for the parameters. Assuming a gravitational wave with a frequency in the kilohertz regime corresponds to a propagation length of a few hundred kilometers during one half period. As in LIGO, this length can be folded into a smaller length scale via retroreflecting mirrors [21]; see Fig. 1. Then, after a propagation time on the order of milliseconds, the light pulses hit the 45° mirrors, which change their direction from $\boldsymbol{K}_{\rm in} = K_x \boldsymbol{e}_x$ to $\boldsymbol{K}_{\rm out} = K_y \boldsymbol{e}_y$, or vice versa. Ideally, this should happen when $\dot{h} = 0$, such that the sign changes of $h(t)$ and the integrand in Eq. (5) cancel each other. Depending on how monochromatic the gravitational wave is, one could repeat this procedure for several half cycles in order to obtain a lasting energy shift of several hE which should then equal or exceed $\hbar\omega$.

In this scheme, the light pulses must fit into one half period of the gravitational wave, such that the pulse duration is well below milliseconds and thus the frequency uncertainty is well above kilohertz. Thus, the idea is to transfer the small energy shift $\Delta E \geq \hbar\omega$ into a phase shift $\Delta\varphi$ which can be measured via interference. To this end, the initial pulse could be split up via a half silvered mirror (a nonpolarizing beam splitter) at 45° into two equal pulses: one first propagating in the x direction and the other one first propagating in the y direction; see Fig. 1. In this way, these two pulses would acquire opposite energy transfers.

Thus far, the setup is similar to a Mach-Zehnder or (half of) a Sagnac interferometer (see also [22–25]) but with an important difference: one does not let the two light pulses interfere at this stage. Instead, they are both sent through another optical path length during which their tiny and opposite energy shifts $\pm\Delta E$ generate a small phase difference $\Delta\varphi$ analogously to a beat note in acoustics. To minimize noise in the relative phase between the two pulses, they can both be sent through basically the same optical path, but in distinguishable optical modes. This phase accumulation period would be after the gravitational wave passed by and thus can be much longer than a period of the gravitational wave. Actually, this fact could be an important advantage in comparison to LIGO, where the effective optical path length $\mathcal{O}(10^3 \text{ km})$ is limited by the period of the gravitational wave such that one does not have more than a few hundred reflections at the mirrors before interference. Of course, this advantage also comes with a drawback since LIGO can measure the full time-dependent amplitude $h(t)$ of the gravitational wave, while the scheme here focuses on the final energy shift.

For photons with energies in the eV regime (visible or near infrared light), a laser pulse with a moderate energy in the millijoule regime contains $N = \mathcal{O}(10^{16})$ photons. In view of their frequency [$\Omega = \mathcal{O}(10^{15} \text{ Hz})$], we see that gravitational waves with $h = \mathcal{O}(10^{-22})$ or even weaker should lead to the stimulated emission (or absorption) of many gravitons with $\omega = \mathcal{O}(\text{kHz})$; cf. [26]. Assuming a classical (coherent) pulse, the usual Poisson limit $\Delta\varphi \propto 1/\sqrt{N}$ yields the achievable phase accuracy $\Delta\varphi = \mathcal{O}(10^{-8})$ for interference measurements. Nonclassical photon states will be discussed below.

Another enhancement factor $\mathcal{O}(10^9)$ is the large ratio of length scales: $\mathcal{O}(\text{km})$ arm length versus $\mathcal{O}(\mu\text{m})$ wavelength. These two enhancement mechanisms are basically the same as in LIGO. As a difference from LIGO, the lasting energy shift ΔE allows longer phase accumulation times. For example, an effective optical path length of $\mathcal{O}(10^6 \text{ km})$, e.g., assuming $\mathcal{O}(10^6)$ reflections (instead of the few hundred reflections at LIGO) would yield a total enhancement factor of $\mathcal{O}(10^{23})$, which looks very promising for amplitudes $h = \mathcal{O}(10^{-22})$.

Form another perspective, each photon acquires a lasting frequency shift of $\pm\Delta\Omega = \mathcal{O}(h\Omega)$ which gives $\mathcal{O}(10^{-7}$ Hz). After a phase accumulation time of a few seconds corresponding to a path length of $\mathcal{O}(10^6$ km), this translates into a phase shift of $\Delta\varphi = \mathcal{O}(10^{-7})$ for each photon—which can then be detected using $N = \mathcal{O}(10^{16})$ photons.

The above considerations assumed that the light pulses are perfectly timed with the gravitational waves such that the former hit the 45° mirrors when $\dot{h} = 0$. If this timing is not perfect, the effect is reduced accordingly. This drawback could be reduced by having a continuous train of (equidistant) pulses such that several pulses are emitted during one gravitational wave period—ideally as a coincidental measurement with LIGO. Note that the total average power of a few watts is not overwhelming. Thus, going to the limit of overlapping pulses, one could also envision a cw laser with a permanent power in this range, where the interference pattern is continuously measured.

Nonclassical photon states—It is well known that one can achieve sensitivities exceeding the Poisson limit $\Delta\varphi \propto 1/\sqrt{N}$ by employing nonclassical states such as squeezed states; see, e.g., [27–29]. Actually, this is being implemented at LIGO; see also [30,31]. Since the energy transfer [Eq. (5)] is bounded by the total energy of the light pulse (independent of its quantum state), a squeezed state would not be an advantage here—except that it could have an energy variance which is different from a coherent state, for example.

However, a nonclassical state can be advantageous for the accuracy of the phase measurement. To understand this point, let us consider the extreme case of a many-body entangled state in the form of a so-called NOON state; see, e.g., [32]. In contrast to a coherent (i.e., classical) state where all photons are in a superposition of the two interferometer arms, this NOON state describes a superposition where either all photons are in one arm (and none in the other) or all photons are in the other arm $|\mathrm{NOON}\rangle = (|N\rangle|0\rangle + |0\rangle|N\rangle)/\sqrt{2}$. After interacting with the gravitational wave, the photons acquire opposite phases $(e^{+iN\Delta\varphi}|N\rangle|0\rangle + e^{-iN\Delta\varphi}|0\rangle|N\rangle)/\sqrt{2}$ such that now the achievable phase sensitivity scales with the Heisenberg limit $\Delta\varphi = \mathcal{O}(1/N)$ instead of the Poisson limit $\Delta\varphi = \mathcal{O}(1/\sqrt{N})$. As a result, the required number N of photons would be much smaller, but actually generating such a highly nonclassical state and measuring it is also much more challenging experimentally.

Quantum aspects of gravity—Thus far, we have assumed that the laws of quantum theory apply to the gravitational field in basically the same way as they do to the electromagnetic field, for example. Now, let us scrutinize this assumption. First, it should be stressed that measuring an energy shift of $\hbar\omega$ does not prove that the energy of gravitational waves is quantized in units of $\hbar\omega$; cf. [33,34]. On the other hand, detecting a gravitational wave at LIGO, for example, and not finding the associated energy transfer in a setup discussed here would indicate that there is something going on that we do not understand (e.g., that the above assumption is wrong).

Furthermore, the setup discussed here could allow us to test certain properties of quantum superposition states of gravitational fields. Similar ideas have already been discussed for the Newtonian gravitational field. If a sufficiently large mass is in a superposition state of two spatially well separated positions, then its (static) gravitational field should also be in a quantum superposition; see, e.g., [35–38]. Going one step further, this superposition state could indicate entanglement between the gravitational field and the matter degrees of freedom—or even mediate entanglement between two different matter degrees of freedom; see, e.g., [39–41].

An analogous idea can be applied to the setup considered here. For example, let us take the NOON state discussed above for the photon field where the light pulse in one arm (say, $|N\rangle|0\rangle$) would gain the energy ΔE while the light pulse in the other arm ($|0\rangle|N\rangle$) would lose this energy. Then, unless one is willing to abandon energy conservation, this means that we get a superposition of quantum states including the gravitational wave, i.e., $|\mathrm{NOON}\rangle|\bar{E}_{\mathrm{grav}}\rangle$ transforms to a superposition of the state $e^{+iN\Delta\varphi}|N\rangle|0\rangle|\bar{E}_{\mathrm{grav}} - \Delta E\rangle$ for one arm and the state $e^{-iN\Delta\varphi}|0\rangle|N\rangle|\bar{E}_{\mathrm{grav}} + \Delta E\rangle$ for the other arm, where $\bar{E}_{\mathrm{grav}}$ denotes the initial energy expectation value of the gravitational wave.

However, this superposition does not necessarily imply strong entanglement between the photon field and the gravitational field, because the quantum states of the latter $|\bar{E}_{\mathrm{grav}} - \Delta E\rangle$ and $|\bar{E}_{\mathrm{grav}} + \Delta E\rangle$ could be nonorthogonal. For example, two coherent states $|\alpha_+\rangle$ and $|\alpha_-\rangle$ have a finite overlap $|\langle\alpha_+|\alpha_-\rangle|^2 = \exp\{-|\alpha_+ - \alpha_-|^2\}$ which can be near unity if the two states $|\alpha_+\rangle$ and $|\alpha_-\rangle$ differ only by an energy of one or a few excitation quanta $\hbar\omega$ on top of a strongly displaced (i.e., nearly classical) state with $|\alpha_\pm| \gg 1$. This would be very different for a Fock state $|n\rangle$, for example, where the overlap $\langle n+1|n-1\rangle$ vanishes (i.e., one has maximum entanglement).

This entanglement between the photon field and the gravitational field or, equivalently, the overlap between the states $|\bar{E}_{\mathrm{grav}} - \Delta E\rangle$ and $|\bar{E}_{\mathrm{grav}} + \Delta E\rangle$ affects the visibility in interference measurements of the phase difference $\Delta\varphi$. For coherent states $|\alpha_\pm\rangle$ of the gravitational field with $\hbar\omega(\alpha_\pm^2 + 1/2) = \bar{E}_{\mathrm{grav}} \pm \Delta E$, the overlap is near unity for small $\Delta E \ll \sqrt{\hbar\omega\bar{E}_{\mathrm{grav}}}$. As a result, the reduced state of the photon field alone is approximately a pure state $(e^{+iN\Delta\varphi}|N\rangle|0\rangle + e^{-iN\Delta\varphi}|0\rangle|N\rangle)/\sqrt{2}$, and thus the phase difference $\Delta\varphi$ can be measured as discussed above; i.e., one has full visibility.

Now let us imagine replacing the coherent state of the gravitational field (in a given mode) by an initial Fock state $|n\rangle$ with $\bar{E}_{\rm grav} = \hbar\omega(n + 1/2)$ and $\Delta E = \hbar\omega$ as an (extreme) example of a nonclassical state. In this case, the overlap $\langle n + 1|n - 1\rangle$ vanishes, and thus the reduced density matrix of the photon field alone is a fully incoherent mixture of the two states $|N\rangle|0\rangle$ and $|0\rangle|N\rangle$, which could in principle be verified by N-photon state tomography. Here, all phase information is lost and one has zero visibility, i.e., the absence of any interference with photonic NOON states. Since a thermal state can be written as an incoherent mixture of (orthogonal) Fock states with their thermal probabilities, the same conclusion (zero visibility) would apply. Thus, by measurements on the photon field (e.g., using the NOON states with variable delay times), one can in principle distinguish between different quantum states of the gravitational field (in this mode), such as between a coherent state and a Fock (or thermal) state (in the absence of other sources of decoherence).

In the bigger picture, the above consideration is an example of gravitational decoherence which has already been discussed in several works; see, e.g., [42–46]. The arguments above are based on the assumption that the laws of quantum theory apply to gravity in the same way that they apply to electromagnetism, for example, but it has also been proposed that one has to modify gravity and/or quantum theory when combining them; see, e.g., [47–49]. In such a case, the predictions could be different (e.g., the decoherence could be larger), and thus the setup could also allow us to test these ideas; see also [50–53].

As another potentially interesting observable, one could measure the phase fluctuations (for the pulsed mode of operation or the cw mode). Since the final phase $\Delta\varphi(t)$ is given by a convolution of the gravitational wave amplitude $h(t')$ with a time-dependent kernel $k(t - t')$ which encodes the history (e.g., reflections) of the photons arriving at a time t, these phase fluctuations $\langle(\Delta\hat{\varphi})^2\rangle$ allow us to access the two-point function $\langle\hat{h}(t)\hat{h}(t')\rangle$ of the graviton field, which also contains information (phase coherence versus thermal fluctuations, etc.) about the quantum state of the graviton field.

Conclusions and outlook—To facilitate the transition from passively observing a natural phenomenon such as a gravitational wave to actively manipulating it, we investigate the stimulated emission or absorption of gravitons by light [54], in analogy to an "optical Weber bar." An important difference to LIGO is the distinction between the interaction time (set by the period and pulse length of the gravitational wave) and the phase accumulation time. For LIGO, both are essentially the same, but in the setup discussed here, the latter is not limited by the gravitational wave but only by optical properties (such as the Q factor) and thus could be much longer. This difference might become even more pronounced for gravitational waves of higher frequencies.

Using nonclassical photon states such as NOON states, energy conservation demands that we create quantum superposition states of gravitational waves with different energies. In this way, interference experiments with variable delay times could even test certain quantum aspects of gravity; e.g., they could distinguish between different quantum states of the gravitational field.

Acknowledgments—Fruitful discussions with S. Liberati, D. Page, and W. G. Unruh are gratefully acknowledged. The author has benefited from the activities of COST Action CA23115: Relativistic Quantum Information, funded by COST (European Cooperation in Science and Technology). Funded by the Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) through the Collaborative Research Center SFB 1242 "Nonequilibrium dynamics of condensed matter in the time domain" (Project No. 278162697). Financial support from and collaborations with the initiative Quantum Simulators for Fundamental Physics (QSimFP) are also gratefully acknowledged. QSimFP is supported by the Science and Technology Facilities Council (Project No. ST/T006900/1) as part of the UKRI Quantum Technologies for Fundamental Physics program.

Data availability—No data were created or analyzed in this study.

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[9] Charged particles such as protons and electrons can also be accelerated to nearly the speed of light, but in this case the electromagnetic radiation would dominate any gravitational effects by far.
[10] Note that this applies to the standard coupling $h^{\mu\nu}T_{\mu\nu}$ to the energy-momentum tensor $T_{\mu\nu}$ in linearized general relativity. For other (hypothetical) coupling terms between pho-

tons and gravitons, the situation might be different. However, even in those cases, symmetries such as Lorentz invariance restrict such photon-graviton conversion processes strongly.

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[53] For theories modifying the standard rules of quantum mechanics in gravity, the resulting experimental implications depend on the specific model under consideration. Furthermore, for each specific model, one should first study whether it does or does not contradict the known experimental results and observations. For example, if one takes the naive semiclassical Einstein equations $R_{\mu\nu} - Rg_{\mu\nu}/2 = 8\pi G_{\rm N}\langle\hat{T}_{\mu\nu}\rangle/c^4$ seriously, one would get averaged gravitational fields for statistical mixtures of a mass at different positions, which is not correct [D. N. Page and C. D. Geilker, Indirect evidence for quantum gravity, *Phys. Rev. Lett.* 47, 979–982 (1981)].

[54] In principle, the stimulated emission or absorption of gravitons could also be achieved via mechanical means, e.g., via a rotating and vibrating barbell; see [16]. However, somewhat similar to a comment in [8], the required parameters make such an experiment quite challenging.

End Matter

Appendix: Wave packets—Let us consider the sequence described above in terms of the wave packets associated with the light pulses. The wave equation obtained from Lagrangian (3) reads

$$(\partial_t^2 - [1-h]\partial_x^2 - [1+h]\partial_y^2)A_z = 0. \tag{A1}$$

Since the frequency $\Omega = \mathcal{O}(10^{15}\ \mathrm{Hz})$ of the electromagnetic waves corresponding to visible or near infrared photons with energies in the eV regime is much larger than the frequency ω of the gravitational wave (e.g., in the kilohertz or hertz range), we may use the WKB approximation. In the coordinates in (2), the wave numbers K_x and K_y are conserved (apart from the reflection at the mirrors), but the frequencies Ω change according to the dispersion relation

$$\Omega^2 = [1-h]K_x^2 + [1+h]K_y^2. \tag{A2}$$

For propagation in either the x or y direction, the energy of each photon thus changes with $\Delta\Omega = \pm h\Omega/2$. During the reflections at the static mirrors (occurring when $\dot{h} = 0$), the frequencies Ω do not change. Hence, by altering the directions as described above, one can transform the momentary changes $\Delta\Omega = \pm h\Omega/2$ into a lasting shift in frequency. Since the total number of photons does not change in this process, we get an energy shift of $\Delta E = \pm hE/2$ for each half period of the gravitational wave, i.e., between two reflections in the mirrors (which is consistent with the above results).

Besides the total energy of the wave packets (on the classical level), let us also consider its shape and amplitude. Since the values of the wave numbers are conserved during free propagation (i.e., between two reflections in the mirrors), the shape of the wave packet does not change in terms of the coordinates in (2). However, owing to the reflections in the mirrors (where the x and y length scales are modified by a gravitational wave), we may get a lasting deformation of the wave packets, i.e., light pulses. Similar to the deflection angle discussed above, these deformations could also be used to detect gravitational waves, at least in principle. However, since these deformations are very small $\mathcal{O}(h)$, we may neglect them in the following and focus on the energy shift. Even after the interaction with the gravitational wave, the energy shift induces a phase difference which grows with the time elapsed and thus can be amplified—while the deformation would not be amplified in the same way.

Finally, for fixed K_x and K_y (i.e., between two reflections in the mirrors), wave equation (A1) simplifies to the ordinary differential equation $\ddot{A}_z + \Omega^2(t)A_z = 0$ with the conserved Wronskian $W = A_z^*\dot{A}_z - \dot{A}_z^*A_z$, which yields $W \approx -2i\Omega|A_z^2|$ in the WKB approximation. Thus, the amplitude of A_z changes with $1/\sqrt{\Omega}$. Since the total energy E of the pulse scales with $(\Omega A_z)^2$ and the volume in terms of the coordinates in (2) does not change during free propagation, we find that E changes proportionally to Ω (i.e., the energy of each photon), as expected. Again, the reflections at the static mirrors (occurring when $\dot{h} = 0$) do not change the total energy—but they transform the Q1 instantaneous changes into a lasting energy shift.