

Nonequilibrium Casimir Force with a Nonzero Chemical Potential for Photons

Kaifeng Chen^{1,2} and Shanhui Fan²

¹*Department of Applied Physics, Stanford University, California 94305, USA*

²*Department of Electrical Engineering, Ginzton Laboratory, Stanford University, Stanford, California 94305, USA*

(Received 12 October 2016; published 22 December 2016)

We introduce a new class of nonequilibrium Casimir forces, where the deviation from equilibrium is achieved through the use of a nonzero chemical potential of photons. Such a force can be observed when two semiconductors are brought in close proximity to each other, and when at least one of the semiconductors is subject to an external voltage. By exact numerical calculations of a sphere-plate configuration, we show that in the total force the non-equilibrium component can dominate over its equilibrium counterpart with a relatively modest external voltage, even when the sphere-plate separation is in the nanoscale. As a result, repulsion can be achieved at the nanoscale even with a relatively modest applied voltage. The results here point to a pathway that can significantly advance the quest for observing and harnessing nonequilibrium Casimir forces in solid-state systems.

DOI: 10.1103/PhysRevLett.117.267401

The study of nonequilibrium Casimir forces represents an emerging direction in the exploration of physical effects induced by quantum and thermal fluctuations [1–9]. In the standard configuration for the study of nonequilibrium Casimir forces, one considers two bodies maintained at different temperatures, and studies the forces that are generated by the quantum and thermal fluctuations in each of the bodies. This is in contrast to the equilibrium Casimir forces that have been extensively studied [10–21] since the pioneering works of Casimir [22] and Lifshitz [23], where the bodies are maintained at the same temperature.

The nonequilibrium situation creates fundamentally different physical effects in fluctuation-induced forces. For example, whereas the equilibrium Casimir forces between two bodies are almost always attractive except with unusual geometries [24,25] or material systems [26,27], it has been shown theoretically that even with regular material systems and geometries the nonequilibrium force can be repulsive, and may provide stable equilibrium points [1,2]. From a practical point of view, having a repulsive interaction is important in nanoelectromechanical systems for preventing the collapse of nanostructures [20,28]. In addition, moving away from equilibrium points to a new dimension in controlling fluctuation-induced forces. For example, in the nonequilibrium case one can control the force by changing the temperature gradient.

In spite of the importance of nonequilibrium Casimir forces for fundamental science and practical applications, observation of such forces has proved to be quite challenging. At small separations, the fluctuation force is dominated by its equilibrium component. To observe a nonequilibrium effect, the distance needs to be at least comparable to the thermal wavelength. At such a distance,

however, the magnitude of the total fluctuation force becomes quite small, which presents a significant experimental difficulty. As a result, the only existing observations of such forces were carried out on atomic systems [9]. There have not been any experiments on nonequilibrium forces in solid-state systems, which are important for practical applications. Moreover, as has been pointed out in Ref. [3], to observe a repulsive interaction between a metallic sphere and a metallic substrate requires a very large temperature difference of more than 600 K; moreover, the repulsion occurs only at a relatively large micron-scale distances. Yet, for nanomechanical applications, it would be far more attractive to be able to achieve repulsive interactions at the nanoscale.

Motivated by the goal to harness the fluctuation-induced forces, in this Letter we propose a class of nonequilibrium Casimir forces that occurs between two bodies possessing different chemical potential for photons. When a semiconductor is under an external voltage bias, the electromagnetic fluctuations of the semiconductor can acquire a nonzero chemical potential proportional to the applied voltage [29–32]. Here we consider two semiconductors subject to different voltages and show that the control of such an external voltage provides a powerful mechanism to harness the fluctuation-induced forces. We show that, in this geometry, the nonequilibrium force exceeds the equilibrium component even when the spacing is at the nanoscale; thus, in this case, the nonequilibrium Casimir force is sufficiently large to facilitate experimental observations. In addition, we show that the forces between the bodies can become repulsive with modest contrast in chemical potentials even at the nanoscale dimension. This represents a capability to achieve repulsion that is beyond what has been possible in standard temperature-based nonequilibrium Casimir forces.

The Casimi-Lifshitz forces arise from the quantum and thermal fluctuations of the system. In a semiconductor under external bias V maintained at temperature T , for the angular frequencies above and below the band gap, the correlations of the fluctuation currents inside the body are assumed to have the following form [31–34]:

$$\langle j_\alpha(\mathbf{r}, \omega) j_\beta^*(\mathbf{r}', \omega') \rangle = \begin{cases} \left(\frac{1}{2} + \frac{1}{e^{[(\hbar\omega - qV)/k_B T] - 1}} \right) \frac{4}{\pi} \hbar\omega^2 \epsilon''(\omega) \delta(\mathbf{r} - \mathbf{r}') \delta(\omega - \omega') \delta_{\alpha\beta} & (\omega \geq \omega_g), \\ \left(\frac{1}{2} + \frac{1}{e^{[(\hbar\omega)/k_B T] - 1}} \right) \frac{4}{\pi} \hbar\omega^2 \epsilon''(\omega) \delta(\mathbf{r} - \mathbf{r}') \delta(\omega - \omega') \delta_{\alpha\beta} & (\omega < \omega_g), \end{cases} \quad (1a)$$

$$(1b)$$

where ω and ω' are the angular frequencies, \hbar is the reduced Planck constant, k_B is the Boltzmann constant, α and β label the directions of polarization, \mathbf{r} and \mathbf{r}' are position vectors, $\delta(\omega - \omega')$ is the Dirac delta function, and ω_g is the bandgap angular frequency. For angular frequencies above the band gap, the current-current correlation is dependent on the external bias V on the semiconductor $p-n$ junction, while for angular frequencies below the band gap, the correlation function reduces to a more commonly used form with $V = 0$, as is shown in Eq. (1b). For such a semiconductor, both Eqs. (1a) and (1b) can be separated into a T -independent term that describes the zero-point quantum fluctuation, and a T -dependent term that describes the thermal fluctuation.

The Casimir force is obtained by calculating the Maxwell stress tensor induced by above-band-gap and below-band-gap fluctuations. From Eq. (1a), by applying the external bias V , one can thus control the contribution to the nonequilibrium Casimir forces from the above-band-gap radiation. To demonstrate the effects created by an external bias, we consider the sphere-plate geometry because such a configuration is commonly used in experiments [35,36] and has been numerically computed by using proximity force approximation [37] or exact calculations using scattering theory [2,3,38–41]. As is shown in Fig. 1(a), the system contains a sphere with radius r and a semi-infinite plate, and the distance from the bottom of the sphere to the plate is d . Throughout the Letter, we denote the sphere as s and the plate as p in super- or subscripts. We denote the temperatures of the sphere, the plate, and the environment as T_s , T_p , and T_{env} , respectively.

In Figs. 1(b) and 1(c), we consider two scenarios of using an external bias to control the nonequilibrium Casimir force. In Fig. 1(b), we consider the case when the plate is made of a semiconductor, and a bias is applied to contacts at P^+ and N^+ regions. In the presence of the bias V_p , using the scattering theory formalism as developed in [3] and noting that the only change in the formalism is the replacement of the equilibrium photon occupation number by a nonequilibrium one, the total force acting on the sphere \mathbf{F}^s is

$$\mathbf{F}^s = \mathbf{F}^{(s,\text{eq})}(T_{\text{env}}) + [\mathbf{F}_p^s(T_p, V_p) - \mathbf{F}_p^s(T_{\text{env}})]^+ + [\mathbf{F}_p^s(T_p) - \mathbf{F}_p^s(T_{\text{env}})]^- + [\mathbf{F}_s^s(T_s) - \mathbf{F}_s^s(T_{\text{env}})]. \quad (2)$$

The first term describes the equilibrium force and the remaining terms describe the nonequilibrium forces. In the remaining terms, the superscript of a force term represents

the body where the force is exerted upon, and the subscript represents the fluctuating sources that generate such force. For example, \mathbf{F}_p^s corresponds to the force on the sphere, due to the fluctuations in the plate. The second term in Eq. (2) describes the contribution to the nonequilibrium Casimir force from the above-band-gap radiation (denoted using a superscript “+”), and is controlled by V_p . The third term in Eq. (2) represents the contribution from the sub-band-gap radiation (denoted using a superscript “-”). The last term describes the contribution from the thermal fluctuations in the sphere because of the difference between T_s and T_{env} . In Fig. 1(c), we consider an alternative case where the sphere is subject to an external bias V_s applied to contacts at its P^+ and N^+ regions. In this case, the total force on the sphere is

$$\begin{aligned} \mathbf{F}^s = & \mathbf{F}^{(s,\text{eq})}(T_{\text{env}}) + [\mathbf{F}_p^s(T_p) - \mathbf{F}_p^s(T_{\text{env}})] \\ & + [\mathbf{F}_s^s(T_s, V_s) - \mathbf{F}_s^s(T_{\text{env}})]^+ + [\mathbf{F}_s^s(T_s) - \mathbf{F}_s^s(T_{\text{env}})]^-. \end{aligned} \quad (3)$$

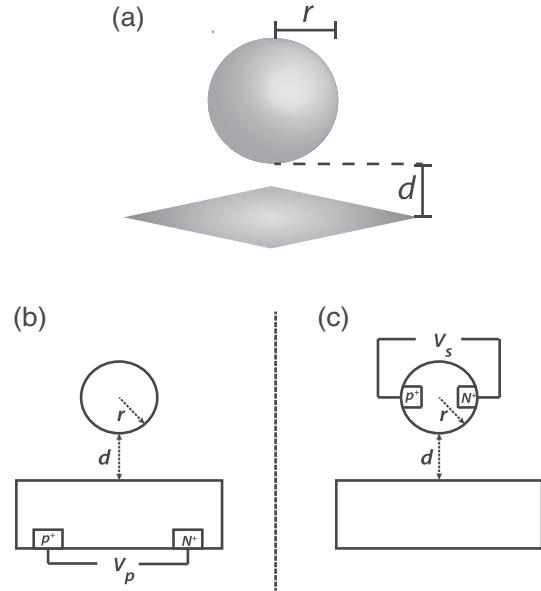


FIG. 1. (a) Configuration of a system consisting of a sphere and a semi-infinite plate. The sphere has radius r and the distance between the bottom of the sphere and the plate is d . (b) An external bias V_p is applied on the p - n junction in the semiconductor plate. (c) An external bias V_s is applied on the p - n junction in the semiconductor sphere.

With $V_p = 0$ or $V_s = 0$, Eq. (2) or Eq. (3) reproduces the standard formula for nonequilibrium Casimir forces, respectively. The explicit form of various terms in Eqs. (2) and (3) are provided in the Supplemental Material [42].

Equations (2) and (3) are applicable to the general nonequilibrium case, where there can be temperature differences for the bodies and the environments, as well as voltage difference between the bodies. For the rest of the Letter, to highlight the effect of the voltage difference, we will consider only the case where $T_s = T_p = T_{\text{env}} \equiv T$. In such a case, Eqs. (2) and (3) are simplified to

$$\mathbf{F}^s = \mathbf{F}^{(s,\text{eq})} + [\mathbf{F}_p^s(V_p) - \mathbf{F}_p^s]^+ \quad (4)$$

and

$$\mathbf{F}^s = \mathbf{F}^{(s,\text{eq})} + [\mathbf{F}_s^s(V_s) - \mathbf{F}_s^s]^+. \quad (5)$$

In Eqs. (4) and (5), as well as for the rest of the Letter, the dependency on T is now implicit because all bodies have the same temperature. Note that, in this case, the only contribution to the nonequilibrium force comes from radiation above the band gap energy.

Using Eqs. (4) and (5), we perform numerically exact calculations on the systems shown in Figs. 1(b) and 1(c), and compute the equilibrium Casimir force and nonequilibrium Casimir force separately. Throughout the Letter, we set $T = 300$ K. For the equilibrium force, i.e., $\mathbf{F}^{(s,\text{eq})}$, following Ref. [39], we perform Wick rotation when integrating over the angular frequencies. The nonequilibrium forces in Eqs. (2) and (3) are obtained using the trace formula introduced in [3] which relates the forces to the traces of the appropriate field correlators. In the numerical calculations, to ensure convergence, we use the criteria given in [43] in the summation over different angular momentum channels.

Because, in the far field, the Casimir force can be viewed as a result of radiative pressure [44], one would prefer to use a material with good light-emitting properties. Moreover, motivated by experiments with sphere-plate geometry that typically measure the forces on the sphere [35,36], we will be primarily interested in the force on the sphere. To ensure absorption of the photons emitted from the plate, we use a semiconductor with a slightly smaller band gap as the material for the sphere. In our calculations, we choose GaAs as the semiconductor for the plate and InP for the sphere. The band gap of GaAs is 1.42 eV at 300 K, while InP has a band gap of 1.34 eV. The dielectric functions for GaAs and InP for the angular frequencies above the band gap are obtained from [45]. For sub-band-gap dielectric function, we use Lorentz-Drude model with the form of $\epsilon(\omega) = \epsilon_\infty(\omega^2 - \omega_L^2 + i\gamma\omega)/(\omega^2 - \omega_T^2 + i\gamma\omega)$ to take into account the contributions from phonon polaritons. For GaAs, ω_T and ω_L correspond to energies of 0.0333 eV and 0.0362 eV, respectively, and $\epsilon_\infty = 11.0$ and $\gamma = 4.52 \times 10^{11}$ rad/s [45]. For InP, ω_T and ω_L correspond

to energies of 0.0377 eV and 0.0428 eV, respectively, and $\epsilon_\infty = 9.61$ and $\gamma = 6.59 \times 10^{11}$ rad/s [45]. The radius of the sphere r is fixed to be 1 μm .

We first consider the configuration in Fig. 1(b). We plot the equilibrium Casimir force $\mathbf{F}^{s,\text{eq}}$, the nonequilibrium Casimir force $\Delta\mathbf{F}_p^s = [\mathbf{F}_p^s(V_p) - \mathbf{F}_p^s]^+$, and the total force on the sphere \mathbf{F}^s in the presence of bias $V_p = 1.36$ V as a function of d from 10 nm to 5 μm in Fig. 2(a) in blue, green, and red curves, respectively. We define the force to be positive when it is attractive. For this system, the equilibrium component, $\mathbf{F}^{(s,\text{eq})}$, is always attractive. At small d , the equilibrium component dominates and the total force is attractive. As d increases, the equilibrium component decreases rapidly. At $d \geq 50$ nm, the nonequilibrium component starts to dominate, and the overall force becomes repulsive. At larger d , in the far-field limit, the total force comes entirely from its nonequilibrium component, which can also be interpreted as radiation pressure [44]; as a result, the total force becomes independent of d . Therefore, in this system, one can observe nonequilibrium force and repulsion at the nanoscale. This is in contrast with previous works on nonequilibrium Casimir force based on temperature gradients, where repulsion is achieved at the micron scale [2].

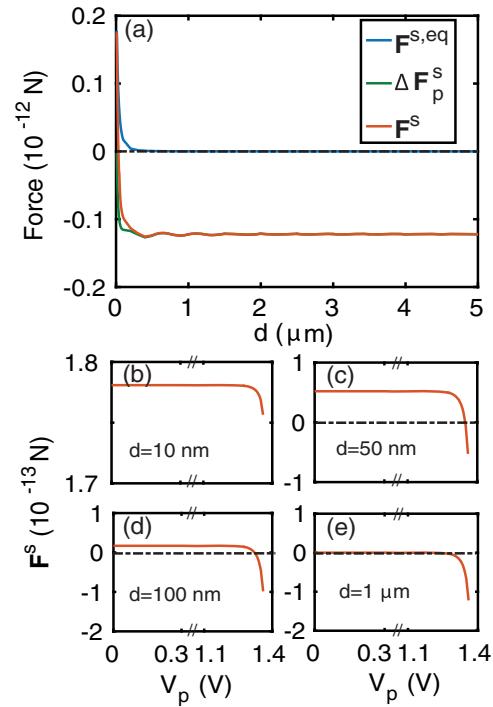


FIG. 2. (a) The equilibrium Casimir force $\mathbf{F}^{s,\text{eq}}$, the nonequilibrium Casimir force $\Delta\mathbf{F}_p^s = [\mathbf{F}_p^s(V_p) - \mathbf{F}_p^s]^+$ and the total force on the sphere \mathbf{F}^s for the configuration of Fig. 1(b) as a function of d from 10 nm to 5 μm when $V_p = 1.36$ V. The dot-dashed horizontal black line corresponds to $|\mathbf{F}^s| = 0$. (b)–(e) The total force on the sphere \mathbf{F}^s as a function of the external bias V_p for various values of d shown in the legend. The dotted-dashed horizontal black line corresponds to $|\mathbf{F}^s| = 0$.

For this system, the nonequilibrium force is strongly influenced by the applied bias V_p . In Fig. 2(c), we plot the total force as a function of applied bias when the spacing is fixed at $d = 50$ nm. When V_p is small compared to the band gap, the total force is dominated by its equilibrium component and, thus, is attractive. However, when the external bias approaches the band gap voltage, the non-equilibrium Casimir force is significantly enhanced and, at $V_p \geq 1.35$ V, it dominates over the equilibrium Casimir force and results in repulsion. We further plot the total force on the sphere as a function of external bias for $d = 10$ nm, $d = 100$ nm, and $d = 1 \mu\text{m}$ in Figs. 2(b), 2(d), and 2(e), respectively. In all these cases, strong deviation of the total force from its equilibrium components occurs only when the voltage is close to the band gap. At $d = 10$ nm, the total force remains attractive for the bias range considered. At $d = 100$ nm and $d = 1 \mu\text{m}$, we observe the transition from attraction to repulsion as the voltage increases; the transition occurs at a smaller voltage for larger d . In this system, repulsions are easily achieved at an external bias that is easily achievable experimentally. This is in contrast with previous works on nonequilibrium Casimir force based on temperature gradients, where a large temperature difference is needed to achieve repulsion [2,3].

The results above are for the configuration of Fig. 1(b), where we apply an external bias on the plate. In contrast, we now instead consider the configuration in Fig. 1(c), where an external voltage is applied to the sphere. In Fig. 3(a), we plot the total forces acting on the sphere as d varies from 10 nm to 5 μm when a voltage of $V_s = 1.26$ V is applied on the sphere. For small separations, the equilibrium component dominates and the total force is attractive. For larger d , we observe an oscillating behavior of the total force. Mathematically, the oscillation arises from the third term in Eq. (3), which describes the self-force on the sphere, i.e., the force on the sphere due to the current fluctuations in the sphere. Physically, such oscillations originate from the interference of the two wave components emitted from the sphere. One of these components is the emission towards the plate, which is then reflected from the plate. The other component is the emission in the direction away from the plate. In contrast to the case in Fig. 2(a), which exhibits only one unstable equilibrium point, here the oscillation results in multiple equilibrium points where the total force is zero. These equilibrium points can be either stable or unstable, as shown in Fig. 3(a). The existence of the stable equilibrium points is a distinct feature of the nonequilibrium force for the configuration of Fig. 1(c). Similar oscillation has been previously noted in the nonequilibrium force generated by a temperature gradient [2]. In our case it is controlled by voltage rather than temperature gradient.

We further examine the dependence of the total force on V_s in Figs. 3(b)–3(e) for four different values of d . For the cases $d = 10$ nm and $d = 20$ nm, as shown in Fig. 3(b), the

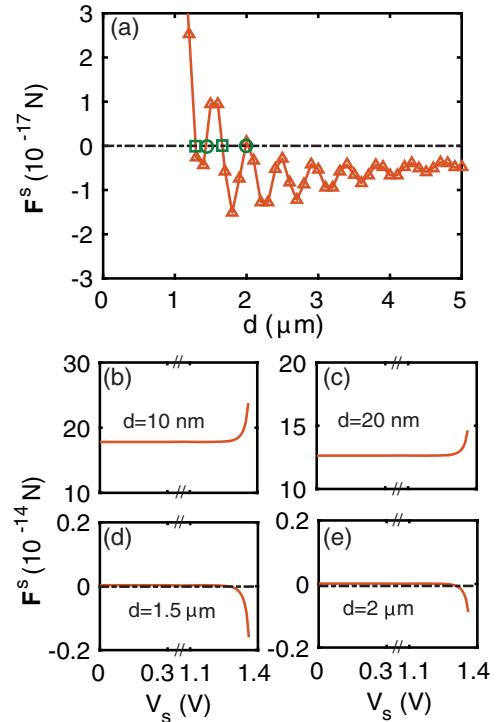


FIG. 3. (a) Total force on the sphere as a function of d for $V_s = 1.26$ V, for case (c) in Fig. 1. The dotted-dashed horizontal black line corresponds to $|F^s| = 0$. Two unstable equilibrium points for are labeled using squares, and two stable equilibrium points are labeled using circles. (b)–(e) The total force on the sphere F^s as a function of the external bias V_s for various values of d shown in the legend. The dotted-dashed horizontal black lines in (d) and (e) correspond to $|F^s| = 0$.

increase of the voltage enhances the attraction of the sphere and the plate. This is in strong contrast with the case considered in Fig. 2. On the other hand, for the cases of $d = 1.5 \mu\text{m}$ and $d = 2 \mu\text{m}$ [Figs. 3(d) and 3(e)], as V_s increases, the total force on the sphere becomes repulsive for large V_s . Because of the interference effects in the self-force of the sphere, this configuration exhibits a rich set of effects in the nonequilibrium forces as we vary either the voltage or the gap size.

In the results above, for ease of exact numerical calculations, we have chosen a relatively small sphere with a radius of 1 μm . Larger force is expected for a larger sphere. In our configuration, for a sphere radius of 5 μm , a repulsive force of 1.93×10^{-12} N can be achieved with a gap size of 50 nm and with a voltage of 1.36 V applied on the plate, as can be estimated using the proximity force approximation that has been shown to be accurate for a sphere of this size in this near-field regime [39]. Forces of such a magnitude are readily detectable with the current experimental techniques for the measurement of Casimir forces [46]. To summarize, we have introduced a new class of nonequilibrium Casimir forces, where the deviation from equilibrium is achieved by applying external voltages to semiconductors to create a nonzero chemical potential for

photons. The use of such a chemical potential should significantly simplify the quest for observing and harnessing nonequilibrium Casimir forces in solid-state systems.

This work was supported by the Department of Energy “Light-Material Interactions in Energy Conversion” Energy Frontier Research Center under Grant No. DE-SC0001293.

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