Giant frequency-selective near-field energy transfer in active–passive structures

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Radiative heat transfer between nearby objects can be much larger in the near field (submicron separations) than in the far field [1–3] due to coupling between evanescent (surface-localized) waves [4,5]. In this paper we investigate the possibility of exploiting both active materials and geometry to enhance and tune near-field energy transfer. In particular, we study amplified spontaneous energy transfer (ASET)—the amplified spontaneous emission (ASE) from a gain medium that is absorbed by a nearby passive object—and demonstrate orders of magnitude enhancements compared to far-field emission or transfer between passive structures. Our work extends previous work on heat transfer between planar, passive media [6–9] to consider the possibility of exploiting both active materials and geometric modifications to the polarization response of spheres (for lattices), revealing not only significant potential enhancements but also strongly geometry-dependent additional enhancement.

Recent approaches to tailoring incoherent emission from nanostructured surfaces have begun to explore situations that deviate from the usual linear and passive materials [12–17], with the majority of these works primarily focusing on ways to control far-field emission, e.g., the lasing properties of active materials [18]. Here we consider a different subset of such systems: structured active–passive bodies that exchange energy among one another more efficiently than they do into the far field. Our predictions below extend recent progress in understanding and tailoring energy exchange between structured materials, which thus far include doped semiconductors [19], phase-change materials [20,21], and metallic gratings [22–24]. Active control of near-field heat exchange offers a growing number of applications, from heat flux control [25,26] and solid-state cooling [26] to thermal diodes [27,28]. Our work extends these recent ideas to situations involving systems undergoing gain-induced amplification.

The starting point of our analysis is the well-known linear fluctuational electrodynamics framework established by Rytov, Polder, and van Hove [29,30]. In particular, given two bodies held at temperatures $T_1$ and $T_2$, and separated by a distance $d$, the power or heat transfer from 1 to 2 is given by [4]

$$P(T_1,T_2) = \int_0^\infty (\Theta(\omega,T_1) - \Theta(\omega,T_2)) \Phi_{12}(\omega) \frac{d\omega}{2\pi},$$

(1)

where $\Theta(\omega,T)$ is the mean energy of a Planck oscillator at frequency $\omega$ and temperature $T$, and $\Phi_{12}(\omega)$ denotes the spectral radiative heat flux, or the absorbed power in object 2 due to spatially incoherent dipole currents in 1. Such an expression is often derived by application of the fluctuation-dissipation theorem (FDT), which relates the spectral density of current fluctuations in the system to dissipation [4], $(J_i(x,\omega), J_j(x',\omega')) = \frac{i}{\pi} \int_0^\infty \frac{\omega \epsilon_0 \Im \epsilon(x,\omega) \delta(x - x') \Theta(\omega,T) \delta_{ij}}{-\omega} d\omega$,
and \( n_2 \), resulting in the following effective permittivity [32]:

\[
\epsilon(\omega) = \epsilon_r(\omega) + \frac{4\pi g^2}{\hbar k^2} \frac{\gamma_1 D_0}{\omega - \omega_{21} + i\gamma_1 e^{i\omega_2/\epsilon}},
\]

where \( \epsilon_r \) denotes the permittivity of the background medium and the second term describes the gain profile \( \epsilon_G \), which depends on the “lasing” frequency \( \omega_{21} \), polarization decay rate \( \gamma_1 \), coupling strength \( g \), and population inversion \( D_0 = n_2 - n_1 \) associated with the \( 2 \rightarrow 1 \) transition. Detailed-balance and thermodynamic considerations lead to a modified version of the FDT \([31,33,34]\) involving an effective Planck distribution \( \Theta(\omega_{21}, T_G) = -n_2\hbar \omega_{21}/D_0 \), in which case the system exhibits a negative effective or “dynamic” temperature under population inversion, a negative effective or “dynamic” temperature under \( \Theta < 0 \) population inversion, i.e., \( \Im \epsilon_G < 0 \), the spectral electric-current correlation function associated with the active medium,

\[
\langle J_j(x, \omega) J_j^*(x', \omega) \rangle = \frac{-4}{\pi} \omega \epsilon_0(\Im \epsilon_G) n_2 \hbar \omega_{21}/D_0 \delta(x - x') \delta_{ij}
\]

is positive. As a consequence, the heat transfer originating from atomic fluctuations in an active body to a passive body always flows from the former to the latter, i.e., \( T < 0 \) reservoirs always transfer energy \([31]\). Of course, in addition to fluctuations of the polarization of the gain atoms, such a medium will also exhibit fluctuations in the polarization of the host medium, depending on its thermodynamic temperature and background loss rate \( \sim \Im \epsilon_r \), as described by the standard FDT \([4]\). Although thermal flux rates can themselves be altered (e.g., enhanced) in the presence of gain through the dependence of \( \Phi_{12} \) on the overall permittivity, the flux rate from such an active medium will tend to be dominated by the fluctuations of the gain atoms, the focus of our work.

I. PLANAR MEDIA

We begin our analysis of ASET by first considering an extensively studied geometry involving two semi-infinite plates that exchange energy in the near field. Such a situation has been thoroughly studied in the past in various contexts \([6–9]\), but with passive materials, whereas below we consider the possibility of optical gain in one of the plates. For simplicity we omit the frequency dependence in the complex dielectric functions \( \epsilon_j \) of the two plates \((j = 1, 2)\), shown schematically in Fig. 1 along with our chosen coordinate convention. We assume that one of the plates is doped with a gain medium, such that \( \epsilon_1 = \epsilon_r + \epsilon_G \), and consider only fluxes due to fluctuations in the active constituents \( \sim \Im \epsilon_G \), as described by the modified FDT above \([4,35]\). Due to the translational symmetry of the system, it is natural to express the heat flux in a Fourier basis of propagating transverse waves \( k_j \) \([4]\), in which case the flux is given by an integral \( \Phi(\omega) = \int \Phi(\omega, k_j) k_j d k_j \). In the near field, \( k_j \gg \omega/c \), the main contributions to the integrand come from evanescent waves which exchange energy at a rate \([5,29]\)

\[
\Phi_{12}(\omega, k_j) \approx \sum_{q_x, q_p} \frac{\Im(\epsilon_G) \Im(r_2^q) \Im(r_1^p) e^{-2\Im|\gamma_0/d|}}{\Im \epsilon_1 |1 - r_1^q r_2^p e^{-2\Im|\gamma_0/d|}|^2},
\]

where \( \epsilon_1 = \epsilon_r + \epsilon_G \), and \( r_1^p = \epsilon_r^{\text{ph}} - \epsilon_r^{\text{ph}} \), are the Fresnel reflection coefficients at the interface between vacuum and the dielectric media, for \( s \) and \( p \) polarizations, respectively, defined in terms of the wave vectors \( k_1 = k_0 \hat{\gamma} + \gamma \hat{z} \), with \( |k_0| = \omega/c \) and \( |k_1|^2 = \gamma^2 + \epsilon_1 \omega^2/c^2 \). Note that the derivation of Fresnel coefficients requires special care since \( \epsilon_1 \approx 0 \), the sign of the perpendicular wave vector \( \gamma_1 = \pm \sqrt{\epsilon_1 \omega^2/c^2 - k_1^2} \) needs to be chosen correctly inside the gain medium \([36–38]\). Here we make the physically motivated choice that yields decaying surface waves inside the semi-infinite gain medium. In the case of evanescent waves \( k_1 \gg \omega/c \), \( \gamma_1 \approx \Im \epsilon_1 \), such that \( r_1^p \rightarrow 0 \) and \( r_1^p = \epsilon_i^{\text{ph}} - \epsilon_2^{\text{ph}} \approx |k_1|^2 + \epsilon_2 \omega^2/c^2 \), where \( \epsilon_i = \epsilon_j + i \epsilon_j' \). Substituting \( e^{2q/d} = z \) and approximating the integral \( \int z f(z) d z \approx z_0 f(z) \), with \( z_0 = k_0 d = \Im(r_1^p r_2^p) \) denoting the wave vector that minimizes the denominator of (4), one obtains

\[
\Phi_{12}(\omega) = \frac{z_0 \Im(\epsilon_G) \Im(r_1^p) \Im(r_2^p)}{4\pi^2 d^2 \Im \epsilon_1} \times \int_1^{\infty} \frac{d z}{\left[z - \Re \left(r_1^p r_2^p\right)^2\right]^2 + \left[\Im \left(r_1^p r_2^p\right)^2\right]^2}.
\]

It follows that the flux rate in the case of passive media with small loss rates scales as \( \Phi_{12} \approx \Im(\epsilon_1) \Im(\epsilon_2) \left|\epsilon_2^{\text{ph}} - \epsilon_1^{\text{ph}}\right|/(4\pi^2 d^2) \sim 1/\ln \left|\epsilon_2^{\text{ph}} - \epsilon_1^{\text{ph}}\right| \) under the resonant condition \( \Re \epsilon_j = -1 \), illustrating a slow, logarithmic dependence on the loss rates and corresponding divergence as \( \Im \epsilon_j \rightarrow 0 \), described in Ref. [11]. However, ASET in the presence of gain, described by (5), depends differently on the loss rates. On the one hand, in situations where gain does not compensate for losses \( \Im \epsilon_i > 0 \), the integral can be further simplified to yield \( \Phi_{12} \approx 1/\Im \epsilon_1 \Im \epsilon_2 \Im \epsilon_3 \), illustrating the same logarithmic dependence on loss rates and resonant conditions, but with the flux rate exhibiting an additional factor \( \sim \Im \epsilon_G / \Im \epsilon_1 \). On the other hand, when the active plate has overall gain, i.e., \( \Im \epsilon_1 < 0 \), the integral diverges under the modified condition \( \Re(r_1^p r_2^p) > 1 \) and \( \Im(r_1^p r_2^p) = 0 \), or alternatively,

\[
(|\epsilon_1|^2 - 1)(|\epsilon_2|^2 - 1) = 4 \epsilon_1''^2 > |\epsilon_1 + 1|^2|\epsilon_2 + 1|^2,
\]

\[
\epsilon_1''^2(1 + |\epsilon_2|^2 - 1) + \epsilon_1''(|\epsilon_2|^2 - 1) = 0,
\]

FIG. 1. Schematic of two semi-infinite plates of permittivities \( \epsilon_1 \) and \( \epsilon_2 \), respectively, separated by a vacuum gap \( d \). Fourier decomposition of scattered waves with respect to parallel \( k_2 \) and perpendicular \( \gamma \) wave vectors simplifies calculations of energy transfer.
both of which cannot be simultaneously satisfied below threshold. Note that in this regime, \( \Re \varepsilon = -1 \) is no longer a necessary condition for maximum heat transfer. In particular, the divergence can occur at unequal values of \( \Re \varepsilon_j \) and \( \Im \varepsilon_j \), in which case the linewidth \( \sim |\Im(r'_1 r'_2)| \) and peak wave vector \( \sim \Re(r'_1 r'_2) \) are decreased and increased, respectively, by suitable choices of material parameters. Such a divergence is of course indicative of a LT, at which point linear fluctuational electrodynamics is no longer valid. Although semi-infinite plates offer analytical insights and computational ease, their close nature and large effective loss rates make them far from ideal for studying ASET. In what follows, we consider finite and open geometries in which even larger ASET and tunability can be attained.

II. SPHERE DIMERS AND LATTICES

A. Sphere dimers

Consider an illustrative open geometry consisting of two spheres separated by vacuum, shown in Fig. 2. In addition to material loss, such a system also suffers from radiative losses, which we quantify (neglecting stimulated emission) from the far-field flux \( \Phi_0 \). The calculation of heat transfer between two spheres was only recently carried out using both semianalytical \([40,41]\) and brute-force methods \([39]\). Here we extend these studies to consider far-field radiation from one of the spheres (in the presence of the other) and the possibility of gain. In particular, we analyze near-field energy exchange between two spheres \( \Phi_{12} \) and far field emission \( \Phi_0 \) by employing a semianalytical method (SA) based on Mie-series expansion of scattered waves, and which follows from a recent study of heat transfer in a similar system by employing field expansions in terms of Mie series \([40]\). Figure 2 shows a schematic of the system, consisting of two vacuum-separated spheres of radii \( R_1 \) and \( R_2 \), dielectric permittivities \( \varepsilon_j \), separated by surface–surface distance \( d \), where one of the spheres is doped with a gain medium, such that \( \varepsilon_1 = \varepsilon_G + \varepsilon_1 \). We compute the flux rates through a surface \( S \) in vacuum from dipoles \( \mathbf{x}_i \in V_1 \) which is given by \( \Re \int_S (\mathbf{E}^i \times \mathbf{H}) = \frac{\omega \Im \varepsilon_G}{\pi} \Im \int S \int d^3 \mathbf{x}_1 \mathbf{G}^\ast \times (\nabla \times \mathbf{G}) \cdot d \mathbf{S} \), where \( \mathbf{G}(\mathbf{x}, \mathbf{x}_i) \) is the Dyadic Green’s function (GF), or the electric field due to a dipole source at \( \mathbf{x}_i \) evaluated at a point \( \mathbf{x} = \mathbf{x}_1 = \mathbf{x}_2 \) in vacuum, with \( \mathbf{x}_j \) denoting the position relative to the center of sphere \( j \), and where we have employed the FDT above to express the flux as a sum of contributions from individual (spatially uncorrelated) dipoles.

When expressed in a basis of Mie modes, the GF from a dipole at a position \( \mathbf{x}_i \in V_1 \) evaluated at \( \mathbf{x} \) is given by \([8]\)

\[
G(\mathbf{x}, \mathbf{x}_i) = i k_0 \sum_{\ell, v = N}^{\ell, v = N} \sum_{m = N} \sum_{q, d = \pm} M_{\ell, v - m}^{(1)q}(k_0 | \mathbf{x}_i),
\]

where \( k_j = \sqrt{\varepsilon_j}/c \), \( \ell \in \mathbb{Z}^+ \), \( |m| \leq \ell \), \( N \) denotes the maximum Mie order, \( C_{\ell q}(q) \) and \( D_{\ell q}(q) \) are standard Mie coefficients \([40,41]\), \( M_{\ell q}(q) \) denote spherical vector waves, \( z_\ell^{(q)} \) and \( z_\ell^{(q)} \) are spherical Bessel \((p = 1) \) and Hankel \((p = 3) \) functions of order \( \ell \), \( z_\ell^{(q)}(x) = \frac{d}{dx} z_\ell^{(q)}(x) \), and \( V^{(p)}_{\ell m} \) are spherical vector harmonics \([42]\).

The advantages of employing spherical vector waves comes from the useful orthogonality relations \([8]\) described in Appendix A, which greatly simplify the calculation of fluxes, requiring integration over \( V_1 \) and over either the surface \( S : |\mathbf{x}_2| \to R_2 \) circumscribing sphere 2 (as derived previously in Ref. \([8]\)) or a far-away surface \( S : |\mathbf{x}| \to \infty \), leading to the following expressions:

\[
\Phi_{12}(\omega) = \frac{R_1 \Im \varepsilon_G}{R_2 \Im \varepsilon_1} \sum_{m, v, q = \pm} \sum_{q, d = \pm} \Im \left( \frac{1}{x_v(R_2)} \right) \Im \left( \frac{1}{x_v(R_1)} \right) \times \left| \frac{z_\ell^{(q)}(k_1 R_1)D_{\ell q}(q)}{z_\ell^{(q)}(k_0 R_2)} \right|^2 |x_v(R_2)|^2,
\]

\[
\Phi_0(\omega) = \frac{2k_0^2 \Im \varepsilon_G}{\pi \Im \varepsilon_1} \sum_{m, v, q = \pm} \sum_{q, d = \pm} y_\ell^{(q)}(R_1) \left( |D_{\ell q}(q)|^2 + |C_{\ell q}(q)|^2 \right),
\]

where \( C_{\ell q}(q) \) and \( D_{\ell q}(q) \) are so-called Mie coefficients \([40]\), \( x_\ell^{(q)}(r) = k_0 r z_\ell^{(q)}(k_0 r) - k_0\ell r z_\ell^{(q)'}(k_0 r) x_\ell^{(q)}(k_0 r) \)

\[
x_\ell^{(q)}(r) = \lim_{R \to \infty} \int \Im \left[ z_\ell^{(q)}(k_0 R) z_\ell^{(q)'}(k_0 R) \right] \times \Im \left[ z_\ell^{(q)}(k_1 r) x_\ell^{(q)}(k_1 r) \right],
\]

\[
x_\ell^{(q)}(r) = \frac{d}{dx} z_\ell^{(q)}(x) \text{ and } y_\ell^{(q)}(r) = \frac{d}{dx} z_\ell^{(q)}(x) \text{, and } k_0 = \omega / c. \text{ We note that } (10) \text{ appears to be new, but we have checked its validity against numerics \([39]\) and also known expressions in the limit } (d \to \infty) \text{ of an isolated sphere } (40). \text{ We also note that the factors of } \Im \varepsilon_1 / \Im \varepsilon_G \text{ in both flux expressions arise because we only consider fluctuations arising from the active constituents (same as in Eqs. (4) and (5) for plates).}

We begin by describing a few of the most relevant radiative features associated with this geometry, focusing on dimers comprising spheres of constant (dispersionless) dielectric permittivities \( \varepsilon_{1,2} \) and equal radii \( R \), which very clearly delineate the operating conditions needed to observe \( \Phi_{12} \gg \Phi_0 \).
assume that one of the spheres (with dielectric $\epsilon_1$) is doped with a gain medium such that $\text{Im} \epsilon_1 < 0$. The top contour in Fig. 3(a) shows $\Phi_0$ from an isolated sphere of $\text{Re} \epsilon = -1.522$ as a function of gain permittivity $\text{Im} \epsilon_1$, illustrating the appearance of Mie resonances and consequently, ASE peaks occurring at $k_0 R \gtrsim 1$. As expected, the LTs (white circles indicate a select few) associated with each resonance occur at those values of gain where (as in the planar case) $\Phi_0 \rightarrow \infty$ and the mode bandwidths $\rightarrow 0$, decreasing with increasing $k_0 R$ (smaller radiative losses). Note that these divergences are obscured in the contour plot by our finite numerical resolution, which sets an upper bound on $\Phi_0$. The middle contour plot in Fig. 3(a) shows that a passive sphere with $\text{Im} \epsilon_2 = 0.05$ in proximity to the gain sphere ($d/R = 0.3$) causes the Mie resonances to couple and split, leading to dramatic changes in the corresponding LTs. Noticeably, while the presence of the lossy sphere introduces additional dissipative channels, in some cases it can nevertheless enhance ASE (decreasing LTs) by suppressing radiative losses [43]. These results are well studied in the literature [18,43] but they are important here because our linear FDT is only valid below LT. Another feature associated with such dimers is the significant enhancement in $\Phi_{12}$ compared to $\Phi_0$ in the subwavelength regime $k_0 R \ll 1$ [44,45], illustrated by the middle/right contours of Fig. 3(a).

Although such near-field enhancements have been studied extensively in the context of passive bodies [4,7,45], as we show here, the introduction of gain can lead to even further enhancements. This is demonstrated by the flux spectra in Fig. 3(b) (corresponding to slices of the contour maps, denoted by white dashed lines), which compare the flux rates of both active (red lines) and passive (blue lines) dimers. The spectra indicate that, while the large radiative components of Mie resonances at intermediate and large frequencies $k_0 R \gtrsim 1$ lead to roughly equal enhancements in $\Phi_{12}$ and $\Phi_0 \sim \Phi_{12}$, the saturating and dominant contribution of evanescent fields and the presence of surface–plasmon resonances in the long wavelength regime cause $\Phi_0 \rightarrow 0$ and $\Phi_{12} \gg 1$ as $\omega \rightarrow 0$. As expected, the existence and coupling of these resonances depend sensitively on $d/R$, occurring at $\text{Re} \epsilon \approx [-2, -1]$ in the limit $d \rightarrow [0, \infty]$ of two semi-infinite plates or isolated spheres, respectively.

B. Dipolar approximation

Since $\Phi_{12} \gg \Phi_0$ in the subwavelength regime, we consider a simple dipolar approximation (DA) [46,47] or quasistatic analysis to understand these enhancements in more detail. In the quasistatic regime, treating the spheres as point dipoles, we find that the flux rates are given by

$$\Phi_{12} = \frac{12 \text{Im} \epsilon_G}{\pi L^5} \text{Im} \alpha_1^{\text{eff}} \text{Im} \alpha_2^{\text{eff}},$$

$$\Phi_0 = \frac{4 \text{Im} \epsilon_G}{\pi \text{Im} \epsilon_1} (k_0 R)^3 \text{Im} \alpha_1^{\text{eff}},$$

where $\alpha^{\text{eff}}$ denote each spheres’ effective anisotropic polarizability (computed by taking into account induced polarization of the dipoles), with parallel ($\parallel$) and perpendicular ($\perp$) components given by [48]

$$\alpha_{1,1/2}^{\text{eff}} = \frac{1 - \frac{\alpha_{1,1/2}}{L}}{1 - \frac{\alpha_0^{\text{eff}}}{L}}$$

$$\alpha_{0,1/2}^{\text{eff}} = \alpha_{1,1/2} + \frac{2\alpha_2^{\text{eff}}}{L} \frac{L^2}{k_0 R},$$

with $\alpha_0 = \frac{3 - 1}{3 + 1}$ denoting the vacuum polarizability of the isolated spheres in units of $4\pi r^3$ and $L = 2 + \frac{d}{\pi}$ their center-to-center distance in units of $R$.

It is well known that in the far-field dipolar limit $d/R \gg 1$, both $\Phi_{12}, \Phi_0 \rightarrow \infty$ under the resonance condition, $\text{Re} \epsilon = -2$ and zero material loss $\text{Im} \epsilon \rightarrow 0$ [10,44,46]. At smaller separations, these two conditions are modified to $|L^5 - \alpha_1(\alpha_2)| = 0$ ($\parallel$ component) or $|L^5 - \alpha_2(\alpha_2)| = 0$ ($\perp$ component) due to changes in the effective polarizability of each sphere. Despite such a modification, in the case of passive dimers, the divergence can only be reached in the limit $\text{Im} \epsilon_1 \rightarrow 0$. For instance, in passive dimers with $\alpha = \alpha_1 = \alpha_2$, $\text{Im} \alpha^{\text{eff}} \rightarrow \infty$ at specific $L^5 = - \text{Re} \alpha$ ($\perp$ component) and $L^6 = \text{Re} \alpha$ ($\parallel$ component) for $\text{Re} \epsilon$ close to $-2$ but only under the condition of zero loss, illustrated in the top contour of Fig. 4(a) for a small $\text{Im} \epsilon_{1,2} = 0.01$. Ultimately, however, the zero-loss quasistatic condition cannot generally be satisfied in finite, passive geometries, resulting in finite flux rates (even in the limit as $\text{Im} \epsilon \rightarrow 0$); essentially, two far-separated ($d \rightarrow \infty$) spheres will not behave as quasistatic dipoles owing to their finite skin depth, except in the limit $R \rightarrow 0$ in which case only
the flux rates per unit volume rather than the absolute rates diverge \[10,49\]. Gain–loss dimers, on the other hand, exhibit diverging flux rates per unit volume rather than the absolute rates. While the DA does not predict such a divergence, which arises due to higher-order scattering effects, it does predict the right scaling of subwavelength resonances (otherwise absent at far-away separations) at \(d \approx 0.317R\) and \(\omega_0 c / \omega R \approx 0.25\) that splits into two resonances at \(d / R \approx 0.306 R\), whose frequencies \(\omega_0^{\pm}\) move farther apart (white dashed lines in the top contour plot) with decreasing \(d\). Such a resonant coupling mechanism results in an ultralarge redshift \(\omega_0^0 \rightarrow 0\) of one of the branches, as \(d \rightarrow d_c\), eventually leading to the quasistatic divergence and better illustrated in the bottom figure of Fig. 4(c), which shows the spectrum corresponding to three different separations, denoted by white dots. While the DA does not predict such a low-d divergence, which arises due to higher-order scattering effects, it does predict the right scaling of \(\Phi_{12}/\Phi_0\) with the various parameters.

The analysis above suggests that a proper combination of gain, geometry, and subwavelength operating conditions can provide optimal conditions for achieving \(\text{ASE} \gg \text{ASE}\) below the LT. In what follows, we consider a more practical and extended geometry, involving lattices of spheres that exchange energy among one another, where one can potentially observe even larger enhancements, leaving open the possibility of further improvements in other geometries \[9,50,51\]. Because exact calculations of flux rates in such a structure are far more complicated \[52\], we restrict ourselves to quasistatic situations that lie within the scope of our DA.

![Image](https://example.com/image.png)

FIG. 4. (a) Flux-transfer rate \(\Phi_{12}\) associated with the sphere dimer system of Fig. 3 under a simple dipolar approximation (DA), in either passive \((\text{Im} \, \epsilon_{1,2} = 0.01, \text{top})\) or active \((\text{Im} \, \epsilon_1 = -\text{Im} \, \epsilon_2 = -0.1, \text{bottom})\) regimes, as a function of \(\text{Re} \, \epsilon_{1,2}\) and \(d/R\). While the flux rate diverges in the active case under total loss compensation, only the rate per unit volume diverges in the case of finite, passive spheres. The validity of the DA for large \(d > R\) is illustrated in (b), which shows also results obtained using the semianalytical (SA) equations \((9)\) and \((10)\). (c) Flux rate spectra \(\Phi(\omega, \alpha)\) (top) and \(\Phi(\omega, \alpha_{PT})\) (bottom) of the dimer system under the \(PT\) symmetry condition, \(\text{Re} \, \epsilon_{1,2} = -1.522\) and \(\text{Im} \, \epsilon_1 = -\text{Im} \, \epsilon_2 = -0.05\), illustrating the splitting of a subwavelength dimer mode as \(d\) changes around a critical \(d_c \approx 0.306 R\). The two branches include both quasistatic \(\omega_0^{(\pm)}\) and subwavelength \(\omega_0^{(\pm)}\) resonances. (d) Flux spectra at three different separations \(d \approx (0.3056, 0.302, 0.3017) R\), marked by the white dots (i), (ii), and (iii), respectively, in the bottom contour in (c).
The combination of reduced loss rates and resonant, near-field enhancements potentially achievable in extended geometries could lead to orders of magnitude larger heat flux rates compared to planar geometries. In fact, as we showed recently in Ref. [11], structures comprising tightly packed, pairwise-additive dipolar radiators can approach the fundamental limits of radiative energy exchange imposed by energy conservation. In what follows, we analyze more realistic versions of such structures, albeit under gain, demonstrating the possibility of achieving significant and widely tunable near-field and material flux enhancements.

We consider two vacuum-separated square lattices of gain–loss nanospheres having equal radii $R$, lattice spacing $t$, and surface–surface separation $d$, depicted in Fig. 5(a). As noted above, the radiation between and from such structures will, to lower order in $(d,t)/R$, depend on the local corrections to the polarizabilities of each individual sphere. The generalization of the DA to consider such a situation yields the following set of equations for the effective polarizabilities of each sphere:

$$\alpha_{\text{eff}} = \frac{1}{(2 + t/R)^3} \sum_{n_1,n_2=0}^{\infty} \frac{n_1^2 + n_2^2 - 2(d/t)^2}{\left[n_1^2 + n_2^2 + (d/t)^2\right]^{3/2}} \alpha_{\text{eff}}$$

in terms of the bare polarizabilities $\alpha_{G,L}^{(0)}$ and structure parameters. (Note that there are three additional equations, which we have chosen to omit, obtained by letting $G \leftrightarrow L$.)

Figure 5 shows (b) $\Phi_{12}$ and (c) $\Phi_0$ in the subwavelength regime $k_0 R = 0.01$, normalized by the dimensionless lattice area $A/R^2 = (2 + t/R)^2$, assuming spheres of $\epsilon_{1,2} = -1.95 \pm 0.01i$ and for various $t = (2,7) R$. To understand the range of validity of the DA with respect to $d/R$, we once again compare its predictions against our semianalytical formulas (SA) in the case of isolated dimers (dotted blue lines), showing excellent agreement in the range $d/R > 1$; note, however, the failure of DA to predict the additional peak at low $d/R \approx 0.2$. Restricting our analysis to large separations, one finds that the presence of additional spheres causes significant enhancements and modifications to the flux rates, leading to complicated, nonmonotonic dependencies on geometric parameters such as $t$. To illustrate the importance of multiple-scattering among many particles, we also show results obtained using a simple pairwise-additive (PA) approximation (dashed lines), in which the flux rates associated with pairs of spheres are individually summed.

Figure 5(d) compares the performance of sphere lattices against that of parallel plates, showing the maximum achievable $|\phi_{12}|/(A/R^2)$ as a function of the relative gain/loss rate $\Im \epsilon_1/\Im \epsilon_2$ for fixed $d/R = t/R = 2$ and multiple loss rates $\Im \epsilon_2 = [0.01,0.1]$ (red and blue lines), varying $\Re \epsilon_{1,2}$ so as to satisfy the resonant condition (obtained and verified numerically). As noted above, whenever $\Im \epsilon_1 < 0$ (loss compensation), it is always possible to choose geometric parameters under which the system undergoes lasing (gray...
shaded region), though this condition can only be obtained analytically for simple structures such as the plates or dipolar spheres above. Below the LT, it is evident that there is significant enhancement in ASET compared to plates, especially as the lattice system approaches the LT. Such an enhancement depends crucially on the loss rates, decreasing with increasing Im $\epsilon_2$, which can be explained by the weak, logarithmic dependence of the planar flux rates on overall loss compensation [11]. Note that as discussed above, at finite $R$, the DA becomes increasingly inaccurate in the limit $\text{Im } \epsilon_1 \to 0$, owing to the finite skin depth effect [10,49]. Our calculations therefore offer only a qualitative understanding of the trade-offs in exploiting particle lattices as opposed to plates. Under losses $\text{Im } \epsilon_2 \approx 0.1$ typical of plasmonic materials, we find that parallel plates exchange more energy compared to sphere lattices for a wide range of gain parameters (except close to the LT), while the latter dominate at smaller Im $\epsilon_2$ and can be greatly enhanced by the presence of even a small amount of gain. Note that while we have chosen to investigate only the case $|\text{Im } \gamma|/|\text{Re } \gamma| \approx 2$ in order to ensure the validity of the DA, potentially larger enhancements are expected to arise at shorter distances or lattice separations, but such an analysis requires a full treatment of ASET in these extended systems, including both finite size and nonlinear effects [35,53]. Nevertheless, our results provide a glimpse of the opportunities for tuning ASET in structured materials.

D. Real materials

The ability to achieve gain at subwavelength frequencies is highly constrained by size and material considerations. In what follows, we describe ASET predictions in a potentially viable material system. Consider a sphere dimer consisting of two ion-doped metallic spheres, shown schematically in the inset of Fig. 6. While there are many material candidates, including various choices of metal-doped oxides and chalcogenides [54], for illustration, we consider a medium consisting of (2 wt. %) Ga-doped zinc oxide (GZO) that is further doped with elements having transition wavelength $\lambda_2 \approx 2.51 \mu m$, and pumped to a population inversion $D_0 = 0.375(h\gamma_1/4\pi \gamma_2^2)$. Also shown is the far-field emission $\Phi_0$ of the isolated gain sphere (green line). The top inset shows the peak ratio $\Phi_0^{\text{max}}/\Phi_0^{\text{min}}$ with respect to changes in $R$, keeping $d/R$ and $D_0$ fixed. (b) Contour plots illustrating variations in $\Phi_0$ (left/middle) and $\Phi_{12}$ (right) with respect to $D_0$, with the black dashed lines indicating operating parameters in (a). (c) Maximum spectral flux rates $|\Phi_{12}(\omega)|R^2/A$ (left) and $|\Phi_0(\omega)|R^2/A$ (right) for extended sphere lattices comprising GZO gain–loss spheres operating at $D_0 = 0.375(h\gamma_1/4\pi \gamma_2^2)$, well below the LT, but of radii $R \approx 0.375/\omega_1$, as a function of $d/R$ and for different values of $t/R$. Also shown are the flux rates of passive lattices (LL, black solid lines), obtained by letting $D_0 = 0$. Corresponding to the first peak occurs slightly above $\text{Im } \epsilon_L \approx 0.37$, which is the threshold gain needed to compensate material loss, at which point Im $\epsilon_1 < 0$. The black dashed lines in the contours denote the operating parameters of Fig. 6(a), confirming that the system lies below the LT. As expected, smaller dimers lead to larger $\Phi_{12} \approx (k_0 R)^3$, as illustrated by the top inset of Fig. 6(a). Figure 6(c) shows the flux rates (red and blue lines) corresponding to extended lattices of spheres comprising the same GZO gain–loss profiles and with radii $R \approx 0.05c/\omega_1 \approx 20 \text{ nm}$ (in the highly subwavelength regime), in a situation where the system is well below the LT,
which occurs at $D_0 = 0.3(\hbar\gamma_\perp/4\pi^2g^2)$. Noticeably, the flux rates are significantly larger than the rates achievable in passive structures (green solid lines).

**III. CONCLUDING REMARKS**

Our achievements shed light on considerations needed to achieve large ASET between structured active–passive materials, attained via a combination of loss compensation in conjunction with near-field effects. While our work follows closely well-known and related ideas in the areas of near-field heat transport and nanoscale lasers (e.g., spasers), the possibility of tuning and enhancing heat among active bodies in the near field is only starting to be explored [26,67]. Our analysis, while motivating and correct in regimes where ASE dominates stimulated emission, ignores important nonlinear and radiative-feedback effects present in gain media as the LT is approached, nor have we considered specific pump and radiative-feedback effects present in gain media as the closely well-known and related ideas in the areas of near-field coupling in thin films, Appl. Phys. Lett. 93, 043109 (2008).

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APPENDIX: VECTOR SPHERICAL HARMONICS

When deriving the flux rates associated with two spheres, we employed the following spherical-vector functions:

$$\mathbf{M}^{(p)}_{lm}(x) = z^{(p)}_{\ell}(kr)\mathbf{V}^{(2)}_{lm}(\theta, \phi), \quad \mathbf{M}^{(p)\perp}_{lm}(x) = z^{(p)}_{\ell}(kr)\mathbf{V}^{(3)}_{lm}(\theta, \phi)$$

where $z^{(p)}_{\ell}$ are spherical Bessel ($p=1$) and Hankel ($p=2$) functions of order $\ell$, $\xi^{(p)}_{\ell}(x) = \frac{1}{\ell+1} \frac{d}{dx} [x z^{(p)}_{\ell}(x)]$, and $\mathbf{V}^{(1)}_{lm}$ and associated spherical vector harmonics [42]

$$\mathbf{V}^{(1)}_{lm}(\theta, \phi) = \mathbf{\hat{r}} Y^{*}_{lm},$$

$$\mathbf{V}^{(2)}_{lm}(\theta, \phi) = \frac{1}{\ell+1} \left( - \phi \frac{\partial Y^{*}_{lm}}{\partial \theta} + i \hat{\theta} \frac{m}{\sin \theta} Y^{*}_{lm} \right),$$

$$\mathbf{V}^{(3)}_{lm}(\theta, \phi) = \frac{1}{\ell+1} \left( \phi \frac{\partial Y^{*}_{lm}}{\partial \theta} + i \hat{\theta} \frac{m}{\sin \theta} Y^{*}_{lm} \right),$$

which satisfy the following orthogonality relations:

$$\int_S \mathbf{V}^{(p)}_{lm} \cdot \mathbf{V}^{(p')\ast}_{lm'} \, d\Omega = \delta_{l'l'} \delta_{pp'} \delta_{mm'},$$

$$\int_S d\Omega \mathbf{V}^{(2)}_{lm} \times \mathbf{V}^{(2)*}_{lm'} \cdot \mathbf{\hat{r}} = - \int_S d\Omega \mathbf{V}^{(2)}_{lm} \times \mathbf{V}^{(3)*}_{lm'} \cdot \mathbf{\hat{r}}$$

$$= - \delta_{l'l'} \delta_{mm'},$$

$$\int_{V_r} dx' \mathbf{M}^{(1)+}_{lm}(x') \cdot \mathbf{M}^{(1)*}_{lm}(x')$$

$$= R^2 \text{Im} \left( k_{\ell} z^{(1)}_{\ell}(k_{R}; k_{\ell}^{(1)*}(k_{R}; k_{\ell}) \right) \frac{\delta_{l'l'} \delta_{mm'}}{k^2_0 \text{Im} \epsilon_{l}},$$

$$\int_{V_r} dx' \mathbf{M}^{(1)-}_{lm}(x') \cdot \mathbf{M}^{(1)*}_{lm}(x')$$

$$= R^2 \text{Im} \left( k_{\ell} z^{(1)}_{\ell}(k_{R}; k_{\ell}^{(1)*}(k_{R}; k_{\ell}) \right) \frac{\delta_{l'l'} \delta_{mm'}}{k^2_0 \text{Im} \epsilon_{l}}.$$

**References**


