

Mie-tronics supermodes and symmetry breaking in nonlocal metasurfaces

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It is usually believed that symmetry breaking in photonic systems leads to weaker confinement, such as in the case of metasurfaces when bound states in the continuum are replaced by quasibound states with lower quality (Q) factors. Here, in contrast, we reveal that symmetry breaking can instead enhance light trapping by strengthening in-plane nonlocal coupling pathways. We consider finite-size arrays of optical resonators supporting Mie resonances (a Mie-tronics platform) and employ diffraction and multiple-scattering analyses to demonstrate that diffractive bands and Mie-tronics supermodes originate from the same underlying Mie resonances but differ fundamentally in their physical nature. Finite arrays exhibit Q -factor enhancement driven by redistributed radiation channels, thus reversing the trends predicted for infinite lattices. We demonstrate that controlled symmetry breaking opens otherwise symmetry-forbidden electromagnetic coupling channels, enabling polarization conversion in nonlocal metasurfaces. These findings establish a unified wave-physics platform linking scattering and diffraction theories. They also outline design principles for multifunctional metasurfaces that exploit nonlocality for advanced light manipulation, computation, and emission.

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I. INTRODUCTION

The field of optical metasurfaces is rapidly evolving from fundamental science to practical technologies, enabling compact platforms for imaging, sensing, energy harvesting, and information processing [1]. Conventional metasurfaces comprise weakly interacting subwavelength elements, allowing a local description. Yet, recent analyses of their limits show that surpassing these constraints demands nonlocal design [2–4]. When interelement coupling becomes strong, collective optical excitations emerge. These excitations result in nonlocal metasurfaces governed by coherent modes extending across the structure. Because these modes are typically excited by plane waves, they can often be interpreted through diffraction theory, earning the name diffractive nonlocal metasurfaces [5]. Understanding such systems benefits from revisiting classical diffraction theory.

In his 1907 “dynamical theory of gratings”, Rayleigh extended Fresnel’s scalar treatment of reflection and diffraction

to periodic metallic surfaces. He noted that Fresnel’s local approximation fails when the grating period approaches the wavelength. In this regime, the recesses of the grating “act as resonators” [6]. Rayleigh thus recognized, decades ahead of modern electromagnetic modeling, that resonance fundamentally alters diffraction and that a “more strictly dynamical theory” is required. Ugo Fano later provided such a framework [7]. He introduced coupling among diffracted and evanescent orders, showing how these “Rayleigh resonators” form quasistationary surface waves. The interference between these surface waves and the continuum gives rise to asymmetric Fano line shapes. The Rayleigh-Fano formalism laid the foundation for modern numerical solvers such as rigorous coupled-wave analysis (RCWA), also known as the Fourier modal method [8]. These tools are used here to compute diffractive bands of nonlocal metasurfaces.

The design of nonlocal metasurfaces often begins with idealized infinite arrays of resonators, where symmetry breaking within the unit cell serves as a versatile tool for tailoring the dispersion and radiation properties of diffractive bands [9–12]. Engineering such diffractive bands has recently fueled advances in analog optical computing and wave-based information processing [13–16]. In practical implementations, however, metasurfaces are finite—a feature that has drawn growing attention [17–21]. While compactness makes them appealing for integration, it simultaneously invalidates the assumption of infinite periodicity inherent to diffraction theory. As the array size decreases, translational symmetry is broken and the momentum conservation underlying Fourier modal methods no longer holds. In finite arrays, collective

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resonances must be treated explicitly. Such resonances are discrete, inherently nonlocal, and strongly governed by the global geometry of the structure. In this regime, Mie-tronics provides a more appropriate physical description.

Mie-tronics has evolved from classical scattering theory into a modern framework for describing light-matter interactions in complex nanophotonic structures. At the microscopic level, collective resonances originate from long-range radiative interactions among Mie modes, namely, localized multipolar resonances of individual dielectric or metallic resonators. These localized resonances form one key element of the Mie-tronics framework [22]. A second key element is the coupling between them, which can be described through translation coefficients connecting the multipolar fields of different resonators [23]. These coefficients quantify how localized resonances interact across the metasurface to form Mie-tronics supermodes. Such supermodes can enhance device performance and enable new design concepts [24–27]. Mie-tronics therefore provides a unified framework for describing both the local response of isolated resonators [28,29] and the collective supermodes of optically coupled resonator arrays.

Mie-tronics supermodes are often dominated by a leading multipolar order while still retaining contributions from higher-order channels. This nonlocal response is wave-mediated: Multipolar displacement-current distributions induced within constituent resonators interact through electromagnetic fields that span many resonators. As a result, the overall system response depends not only on the properties of individual resonators, but also on collective phase relations and excitation conditions. This collective coupling produces strong spatial dispersion beyond weakly nonlocal homogenized descriptions [30]. Such coupled-resonator arrays are commonly described as resonant metasurfaces supporting collective Mie resonances [31] or surface lattice resonances [32,33]. In this work, we refer to them as nonlocal metasurfaces to emphasize the lateral optical communication among constituent resonators, which serves as a key design parameter for advancing flat-optics applications [3]. This emphasis aligns with the broader view that lateral transport of light within a structure helps determine the capabilities and fundamental limits of flat-optics devices [34]. Understanding how to control such lateral communication in realistic finite metasurfaces is therefore essential for pushing nonlocal flat optics toward higher performance.

Building on this motivation, we systematically investigate the effects of symmetry breaking in nonlocal metasurfaces from two complementary perspectives: diffractive bands of idealized infinite arrays and Mie-tronics supermodes of finite arrays. Contrary to expectations based on the diffractive bands of infinite arrays, symmetry breaking in finite metasurfaces can enhance in-plane multiple scattering and increase the Q factor of certain nonlocal supermodes. We further demonstrate that coupling among Mie modes produces additional bonding and antibonding supermodes beyond the Fourier-mode couplings of classical diffraction theory. This effect expands the functional landscape of nonlocal metasurfaces. Notably, antibonding supermodes remain robust in the presence of a quartz substrate, whereas bonding supermodes are suppressed. Finally, we show that distinct unit-cell geometries support identical Mie-resonant supermodes. This result

highlights the robustness and design versatility of the Mie-tronics framework for both diffractive and finite-size nonlocal metasurfaces.

II. FUNDAMENTALS OF MIE-TRONICS

This section introduces the key concepts required to describe light interaction with individual particles and particle arrays. We first review the historical origins of multipole scattering theory, and then clarify the distinction between diffraction- and scattering-based descriptions. Finally, we introduce Mie modes and multipole coupling as the building blocks of a design framework for nonlocal Mie-tronics systems.

A. Historical origins of Mie-tronics

The intellectual roots of Mie-tronics trace back to the 19th-century work of Clebsch. At a time when lens and mirror design was largely governed by geometrical principles, he sought a rigorous wave-based framework for optical devices [35]. Using both plane-wave and pointlike sources, he introduced an early form of the multipole expansion to solve boundary conditions—an approach that later became a cornerstone of scattering theory. Although Clebsch ultimately declared his effort a failure, his pioneering use of the multipole expansion has become a powerful tool for describing wave interactions with structured media [36]. What once seemed unattainable is now well established: The laws of reflection and refraction follow directly from wave physics [37]. In particular, these laws can be derived from the scattering coefficients of spherical boundaries [38,39]. In this sense, Clebsch’s attempt to reinterpret bulky optical elements in wave terms resonates with today’s drive to realize ultrathin metasurfaces that harness wave phenomena to replace conventional optics.

Clebsch’s ideas influenced Lorenz, who first derived the exact solution for plane-wave scattering by a sphere. This result, however, later became more widely associated with Mie because his systematic multipole analysis provided the analytical foundation for modern scattering theory [40]. Unlike Lorenz, Mie was primarily inspired by Rayleigh, who had independently investigated acoustic and optical scattering problems similar to those considered by Clebsch. Notably, Mie introduced the now-standard term “Rayleigh scattering” to distinguish the point-particle regime from his treatment of larger spheres [41].

Over the following century, these contributions coalesced into the framework of generalized Lorenz-Mie theory, emphasizing both the intrinsic Mie modes of scatterers and the multipole content of the excitation sources [42]. Advances in nanoscience have extended scattering analysis to particles of arbitrary shape [43], enabling new classes of photonic devices [44]. This convergence of multipole physics and nanophotonics defines the emerging paradigm of Mie-tronics [45,46]—a framework that unifies classical scattering theory with modern metasurface design.

B. Diffraction and scattering

The early development of scattering theory was closely intertwined with the study of diffraction. In the 19th century,

Clebsch, Lorenz, and Rayleigh often described light-particle interactions in the language of diffraction [47]. Before Maxwell's electromagnetic formulation, when the Huygens-Fresnel principle dominated, the terms “diffraction” and “scattering” were used almost interchangeably. With the establishment of Maxwell's equations, however, diffraction came to be recognized as a specific manifestation of the broader phenomenon of electromagnetic scattering [48,49]. In the modern electromagnetic picture, an incident field induces multipolar polarization currents in a particle. These induced currents act as secondary sources that reradiate electromagnetic waves, forming what is generally called the scattered field. In specific geometries, the same electromagnetic response may appear as reflected, refracted, or diffracted waves, depending on the structure, boundary conditions, and excitation configuration.

Nevertheless, even today, the terms “diffraction” and “scattering” persist in distinct theoretical contexts, sometimes leading to ambiguity. For example, in diffractive nonlocal metasurfaces, “scattering” is often used to describe the radiative leakage of bound states defined within idealized diffraction theory [31]. For metasurfaces with subwavelength periods, classical diffraction theory typically treats the zeroth Fourier order as the only open radiation channel. However, electromagnetic waves support multiple scattering pathways because of their positive-definite energy spectrum [50]. This fundamental property makes trapping light more difficult than trapping electrons, which is readily achieved in negative potentials, such as quantum wells [51]. In Mie-scattering systems, strong confinement therefore requires high-index particles that support multipolar resonances. These additional channels, associated with multipolar displacement currents induced in the scatterers, extend beyond the standard Fourier harmonics [52]. Accounting for them is useful for engineering superlattices with tailored spectral responses [53]. In this multipolar scattering picture, high- Q nonlocal supermodes arise from constructive interference among several multipole radiation channels rather than from suppressing scattering itself.

For conceptual clarity, we adopt the following convention. Results obtained using RCWA are referred to as diffraction theory. In this approach, electromagnetic fields are expanded into Fourier harmonics under plane-wave excitation and in-plane periodic boundary conditions. This formulation originates from the classical grating analyses of Rayleigh and Fano [6,7]. In contrast, results derived from multipole expansions of fields scattered by the metasurface under dipole excitation are referred to as scattering theory. These calculations are performed under open (radiative) boundary conditions. This framework is rooted in the pioneering works of Clebsch, Lorenz, Rayleigh, and Mie [35].

As shown below, the Mie-mode basis provides a clear connection between diffraction and scattering theories. In this unified picture, Mie-tronics bridges Fourier-mode analyses with multipole expansions. It reveals their common foundations and provides a systematic approach for designing nonlocal multipole-based metasurfaces.

C. Mie modes as the building blocks

Although practical scattering systems are more complex than a single sphere, Mie coefficients and multipole

expansions remain fundamental tools for understanding their optical response. We therefore briefly review the standard Mie formalism, its implementation, and physical meaning. In this framework, the incident field on a sphere $\mathbf{E}^{\text{inc}}$, the field scattered by the sphere $\mathbf{E}^{\text{sct}}$, and the internal field inside the sphere $\mathbf{E}^{\text{int}}$ are expanded in a multipole basis:

$$\mathbf{E}^{\text{inc}}(\mathbf{r}) = \sum_{l=1}^{L_{\text{inc}}} \sum_{m=-l}^l [p_{lm} \mathbf{N}_{lm}(k\mathbf{r}) + q_{lm} \mathbf{M}_{lm}(k\mathbf{r})],$$

$$\mathbf{E}^{\text{sct}}(\mathbf{r}) = \sum_{l=1}^{L_R} \sum_{m=-l}^l [a_l p_{lm} \mathbf{N}_{lm}^{(1)}(k\mathbf{r}) + b_l q_{lm} \mathbf{M}_{lm}^{(1)}(k\mathbf{r})],$$

$$\mathbf{E}^{\text{int}}(\mathbf{r}) = \sum_{l=1}^{L_R} \sum_{m=-l}^l [c_l p_{lm} \mathbf{N}_{lm}(n_s k\mathbf{r}) + d_l q_{lm} \mathbf{M}_{lm}(n_s k\mathbf{r})],$$

where a_l and b_l (c_l and d_l) are the external (internal) Mie coefficients, determining how each multipole component radiates into free space or induces displacement currents within the sphere. Here, n_s denotes the refractive index of the sphere and $k = 2\pi/\lambda$ the free-space wave number.

When the incident field has an analytical multipole representation with order L_{inc} , the solution is fully analytical. In the original Mie formulation, however, the series extends to infinity and must be truncated at a finite order L_R , yielding a semianalytical solution. An empirical rule for a sphere of radius R is [54]

$$L_R \approx 2\pi R/\lambda + 4.05\sqrt{2\pi R/\lambda} + 2. \quad (1)$$

For example, a sphere with the parameters listed in Table I requires L_R between 7 and 10 for wavelengths near 1500 nm. While our focus is on magnetic dipole (MD, $l = 1$) interactions, setting $L_R = 1$ is inadequate because it neglects higher-order coupling, leading to inaccurate estimates of Q factors and field enhancements. Moreover, when $L_{\text{inc}} > L_R$, the higher-order incident components ($l > L_R$) cannot couple to internal Mie modes and thus propagate unscattered.

Mie coefficients are formally valid only for isotropic, homogeneous spheres. For arbitrary geometries—such as the hole, square, and T-shaped unit cells in Fig. 1—full-wave simulations are required. We therefore employ the commercial software Lumerical FDTD to evaluate the Q and Purcell factors. The Purcell factor is defined as P/P_0 , where P and P_0 denote the power dissipated by a dipole source in the presence of the metasurface and in the corresponding homogeneous medium, respectively. The Q factor is defined as $\lambda/\Delta\lambda$, where λ is the resonance wavelength and $\Delta\lambda$ is the full width at half maximum of the spectral response. The reflectance, computed using Lumerical RCWA, is defined as the ratio of reflected to incident power integrated over a unit cell. Further details on their formal definitions and derivations are provided in our previous studies [31,54]. All FDTD project files, simulation setups, and MATLAB scripts used in this work are publicly available on Zenodo [55].

Figure 1(a) shows a schematic of a photonic-crystal slab interacting with an x -oriented MD emitter. This system can produce a vortex beam with suppressed radiation in the normal direction, a key feature discussed in this work. The suppression originates from destructive interference between

TABLE I. Geometrical parameters of the unit cells shown in Fig. 1.

Unit cell	Period (nm)	Diameter/width (nm)	Thickness (nm)	Material	Environment
Sphere	425	$D_s = 420$	—	Si ($n = 3.5$)	Vacuum
Hole	529	$D_h = 350$	$H = 600$	Si slab ($n = 3.5$)	Vacuum
Square	720	$W = L = 400$	$T = 500$	Si ($n = 3.5$)	Vacuum
T-shape	720	$W_1, L_1, W_2, L_2 = 400, 240, 280, 160$	$T = 500$	Si ($n = 3.5$)	Vacuum

multipolar displacement currents induced in the system. These multipolar currents collectively form an eigenmode whose fields are symmetry mismatched with the normal radiation channel. Within diffraction theory, this eigenmode is commonly identified as a symmetry-protected bound state in the continuum (BIC) [31]. In Mie-tronics, this BIC manifests as a vortex beam [56].

From a Mie-tronics perspective, the unit cell shown in the right inset of Fig. 1(a) is functionally equivalent to the three other unit cells in Fig. 1(b). Arrays of these geometries are adapted from previous studies [56–58], with their parameters summarized in Table I. Despite differences in size and shape, we show that these arrays support the same collective MD resonances. Geometry transformations and symmetry breaking of the unit cells, however, modify the electric and magnetic

field distributions of the Mie modes and open additional coupling pathways for controlling light-matter interactions.

Before addressing other geometries, we first examine the spectral profiles of the Mie coefficients and how supermodes (collective resonances) arise in finite arrays of spheres, as well as their connection to Bloch-band formation. Figure 2 shows the spectral profiles of the two lowest-order Mie coefficients: the electric (ED) and magnetic dipoles. The internal ED and MD resonances exhibit broad spectral distributions. For a sphere with unit index contrast, the internal coefficients equal 1; thus, values exceeding unity indicate enhanced energy storage. The ED mode spans nearly the entire wavelength range shown, while the MD mode extends from 1200 to 1940 nm.

This broad spectral overlap explains why single-mode models, such as the dipole approximation, cannot accurately capture the Q and Purcell factors of optically coupled arrays. In these systems, multiple multipoles contribute to the scattering process that governs wave behavior [58]. Their contributions are set partially by their intrinsic spectral profiles. The role of spectral overlap in Mie-tronics is discussed further in Sec. II E.

Despite their broad spectra, the ED and MD resonances correspond to whispering-gallery modes, consistent with both geometrical and wave-optics descriptions. Early photonic-crystal studies argued that large unit cells would not support photonic bands due to the dominance of geometrical optics [51]. Mie-tronics revises this view: All Mie modes inherently embody wave physics and can couple collectively. As shown in prior work, coupling between low- and high-order modes gives rise to photonic bands [54,59].

Next, we show that coupling these intrinsic Mie modes results in delocalized collective modes—Mie-tronics supermodes—in finite arrays. We then examine how these supermodes connect to Bloch bands in infinite periodic systems.

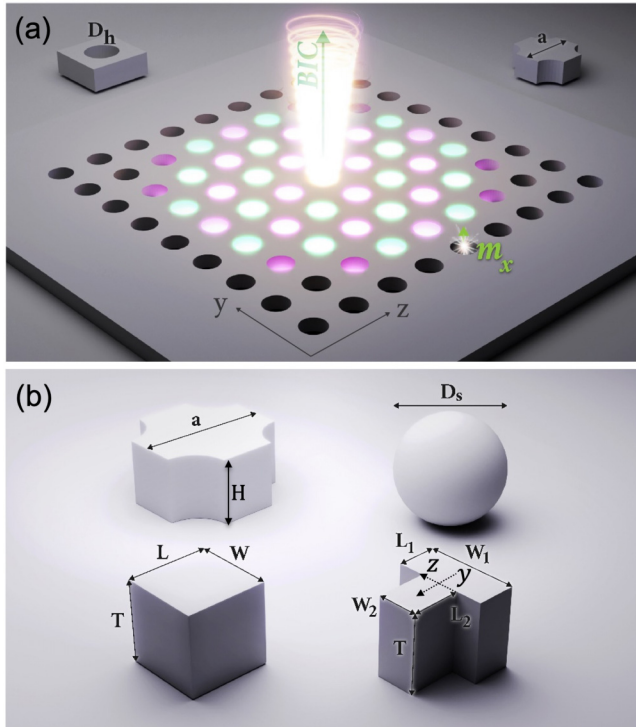


FIG. 1. Beyond spheres: functionally equivalent unit cells in the Mie-tronics framework. (a) Schematic of a magnetic dipole (m_x) interacting with an array of air holes in a silicon slab. The insets show two unit cells that are equivalent from a photonic-crystal perspective but distinct in the Mie-tronics description. (b) Four functionally equivalent unit cells with hole, sphere, square, and T-shaped geometries. Symmetry breaking of the T-shaped unit cell along the z axis redistributes radiation channels, enabling additional coupling pathways.

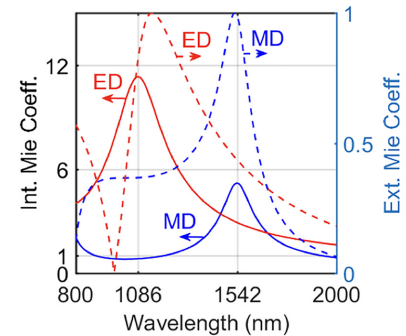


FIG. 2. Spectral dependence of the lowest-order Mie coefficients of the silicon sphere listed in Table I. Internal (left) and external (right) ED and MD coefficients reveal broad overlapping resonances.

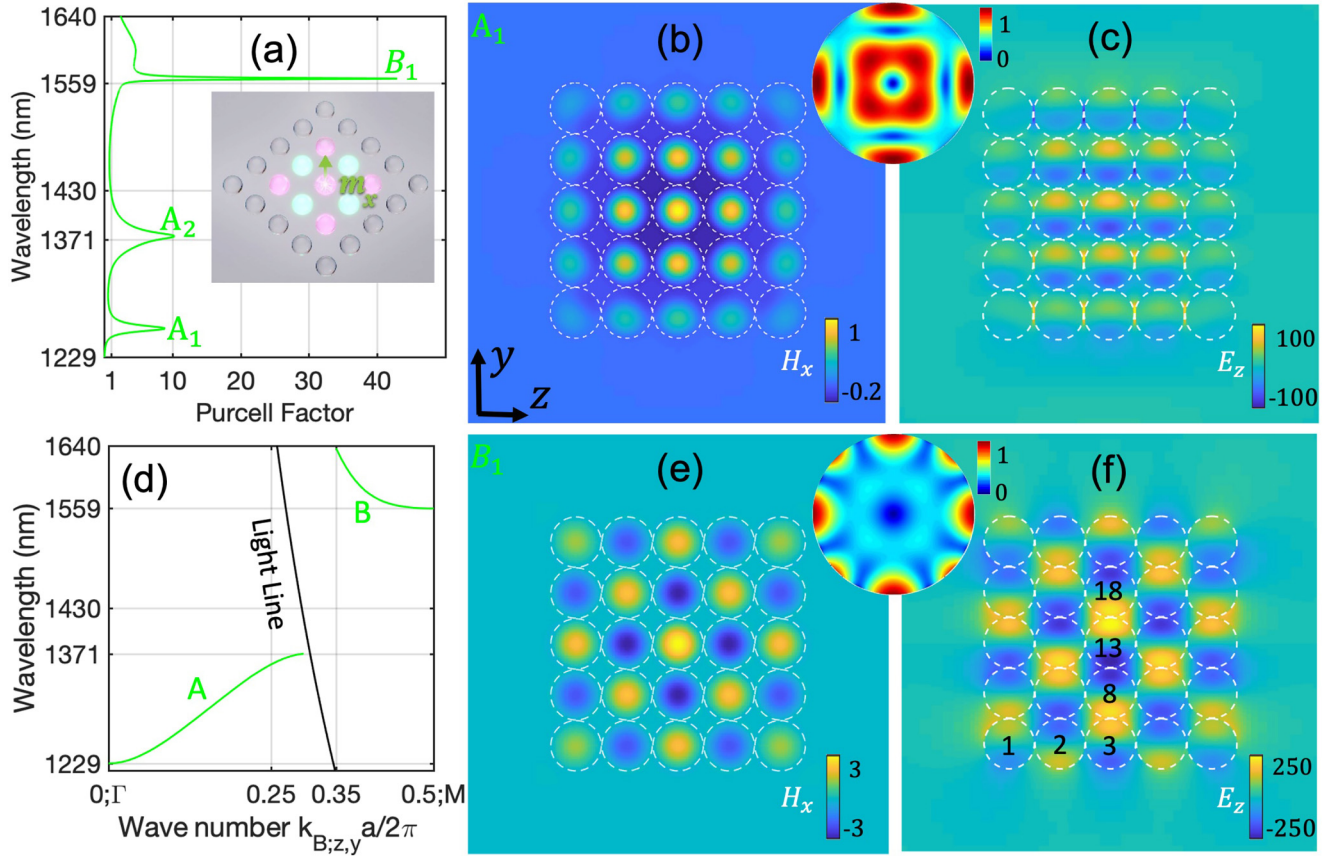


FIG. 3. Mie-tronics origin of supermodes in finite arrays and their correspondence to Bloch bands of the infinite photonic crystal. (a) Purcell-factor spectrum for the configuration in the inset, revealing two well-separated families of collective resonances (antibonding and bonding) with pronounced supermodes labeled $A_{1,2}$ and B_1 . (b), (c) Magnetic (H_x) and electric (E_z) field distributions of the antibonding supermode A_1 ; the inset shows the corresponding far-field radiation pattern. (d) Bloch bands of the corresponding infinite photonic crystal, showing the bands labeled A and B associated with the antibonding and bonding supermodes in panel (a). The vertical axis is plotted in wavelength; accordingly, the bands A and B lie above and below the light line when interpreted in the conventional frequency-based representation. (e), (f) Same as panels (b) and (c), but for the bonding supermode B_1 . The strong electric-field hot spot in the gap between the 13th and 18th spheres indicates efficient coupling of this collective MD mode to a z -oriented ED source.

D. Beyond Bloch bands: Antibonding and bonding supermodes

Figure 3(a) presents the Purcell-factor spectrum for the configuration shown in the inset, where an MD source m_x couples to a 5×5 array of spheres. Two distinct families of collective resonances appear on either side of the peak wavelength (1542 nm) of the intrinsic MD mode shown in Fig. 2. Each family contains well-defined supermodes labeled A_1 , A_2 , and B_1 . The A_1 and B_1 supermodes correspond to the Γ - and M -point edges of the antibonding (A) and bonding (B) Bloch bands in Fig. 3(d), which emerge in the limit of an infinite array. In the photonic-crystal framework, A_1 and B_1 correspond to a guided resonance and a guided mode, respectively. To obtain the dispersions in Fig. 3(d), the dipole m_x is placed at the sphere center, and Bloch boundary conditions are applied along the y and z directions.

Because the multipole modes form an orthogonal basis, the source initially excites only the MD resonance. For this mode, the internal field at the sphere center is determined by the induced MD moment $\eta_x = \sqrt{\frac{3}{8\pi}} \lambda H_x$, where H_x is the local magnetic field [60]. The field in each unit cell,

however, generally contains contributions from multiple multipoles due to interparticle coupling. Linton *et al.* combined multipole expansions with Bloch boundary conditions to describe such multipolar interactions in periodic lattices [61]. This semianalytical approach explains the quantitative agreement between photonic-crystal band-edge simulations and Mie-tronics supermodes in sufficiently large arrays [59]. This agreement holds even though Mie-tronics is formulated in three-dimensional, fully open systems with an infinite number of radiation channels.

The antibonding and bonding characters of these supermodes are evident from their near- and far-field profiles. Figures 3(b) and 3(c) show the A_1 (antibonding) supermode, while Figs. 3(e) and 3(f) show the B_1 (bonding) supermode. The in-plane electric-field (E_z) distributions reveal that A_1 exhibits nodes between adjacent spheres, consistent with in-phase oscillation of the fields [Fig. 3(b)]. In contrast, B_1 displays pronounced hot spots within the interparticle gaps [e.g., between the 13th and 18th spheres in Fig. 3(f)], corresponding to out-of-phase oscillations in neighboring spheres [Fig. 3(e)]. These enhanced local fields enable strong coupling

to emitters or detectors oriented along the z direction. Such collective modes therefore provide a platform for mediating optical interactions between integrated transmitters and receivers and may enable photonic networks for emulating interacting physical systems [62–65].

A further distinction arises from the magnetic-field distributions H_x in Figs. 3(b) and 3(e), which are proportional to the MD moments induced in the spheres [60]. For A_1 , the magnetic moments are aligned, whereas in B_1 they are antialigned. As a result, B_1 extends its fields further into the surrounding environment, while A_1 confines its fields more tightly within the unit cells. This difference in field localization explains the greater robustness of the antibonding supermode to perturbations introduced by a quartz substrate, as discussed in Sec. III.

The far-field radiation patterns [insets of Figs. 3(b) and 3(e)] further clarify the nature of the two supermodes. The bonding supermode B_1 radiates mainly along directions parallel to the metasurface, consistent with its guided-mode character. In contrast, the antibonding supermode A_1 , as a guided resonance, radiates strongly into oblique directions away from the surface. Similar to the hole arrays discussed in the accompanying manuscript [66], the A_1 supermode produces a vortex beam with suppressed radiation along the surface normal, confirming its collective MD origin. From an applications perspective, the in-plane confinement of B_1 makes it especially appealing for on-chip photonic devices that demand minimal radiation losses away from the plane.

Interestingly, neither A_1 nor B_1 radiates along the surface normal (the x axis), preventing excitation by normally incident plane waves. This symmetry protection follows directly from their near-field profiles. A normally incident z -polarized plane wave excites only supermodes with E_z even in z , whereas both A_1 and B_1 exhibit odd symmetry. Similarly, a y -polarized plane wave couples only to supermodes with E_y even in y , but here E_y is odd, again forbidding coupling. These symmetry constraints can equivalently be interpreted in terms of the parity of the Bloch eigenfields within the unit cell. Such symmetry considerations provide useful guidelines for tailoring diffractive bands through controlled symmetry breaking, as demonstrated in Sec. IV. Nevertheless, the collective nature of supermodes allows normal-incidence excitation in sufficiently large arrays. In Fig. 3(a), only three discrete supermodes appear for the 5×5 array, but their number increases rapidly with array size. Because of retardation effects, some of these additional supermodes exhibit E_z distributions along the z axis, enabling coupling to normally incident plane waves [31].

The discrete peaks in Fig. 3(a) arise from distributed scattering with radiative boundary conditions across the finite array. They correspond to high- Q collective supermodes that are not captured by conventional Bloch-wave analysis. In this framework, Bloch waves propagate without scattering in an infinite lattice and form continuous bonding and antibonding bands [Fig. 3(d)]. In finite arrays, edge reflections can confine these waves and produce Fabry-Pérot-type modes. However, the Mie-tronics supermodes originate from distributed multiple scattering rather than edge reflections. The scattered fields emerge predominantly from the interior of the array, not from the edges, as expected for Fabry-Pérot feedback [59]. Distinguishing these mechanisms is nontrivial and underpins unconventional light localization with Mie-tronics [66].

E. Strong nonlocality in Mie-tronics

Nonlocality appears in several related but distinct contexts in electromagnetism. Even free-space propagation is spatially nonlocal in an input-output sense, since diffraction causes the field at a given observation point to depend on an extended region of the preceding wave front. In electromagnetic metamaterials, however, nonlocality is more commonly discussed as spatial dispersion, namely, a wave-vector-dependent material or effective-medium response [30]. Recent progress in photonic crystals, metasurfaces, and flat optics has further emphasized nonlocality as an enabling design degree of freedom, rather than merely a correction to local response models [3]. Across these areas, multipole expansion provides a powerful framework for connecting induced polarization currents within constituent elements to the macroscopic electromagnetic response, thereby guiding structural design and optimization.

For clarity, this subsection distinguishes how multipole expansion is used in two regimes. The first is weak nonlocality, or weak spatial dispersion, where an effectively continuous description may still be meaningful. The second is strong nonlocality, or strong spatial dispersion, where a unique set of effective material parameters generally becomes ill-defined. To facilitate a qualitative discussion, we follow the simplified formulation of Ref. [30]. In this formulation, spatial nonlocality is expressed through the dependence of the polarization current density at an observation point $\mathbf{r}$ on the electric field over an extended spatial region,

$$\mathbf{J}(\mathbf{r}) = \int_{V_\infty} K(\mathbf{r} - \mathbf{r}') \mathbf{E}(\mathbf{r}') d^3 \mathbf{r}',$$

where K is a spatial response kernel. This expression shows that the response at $\mathbf{r}$ is not determined solely by the field at the same point, but by the field distribution in its surrounding region V_∞ . Equivalently, in Fourier space, such a nonlocal relation corresponds to a wave-vector-dependent response.

When the medium can still be regarded as effectively continuous and the spatial dispersion is weak, the nonlocal response can be expanded perturbatively in spatial gradients of the electric field. The polarization current density can then be written as

$$\mathbf{J}(\mathbf{r}) = j\omega\kappa\mathbf{E}(\mathbf{r}) + j\omega\kappa'\nabla\mathbf{E}(\mathbf{r}) + j\omega\kappa''\nabla\nabla\mathbf{E}(\mathbf{r}) + \dots$$

Here, κ , κ' , κ'' , $\dots$ represent effective material parameters, with tensor contractions implied. The first term corresponds to the local ED response, whereas the higher-gradient terms encode the leading effects of spatial dispersion. These field-gradient terms are commonly associated with multipolar contributions beyond the ED response, such as MD and electric-quadrupole responses, which are treated as part of the intrinsic effective material response in the macroscopic description [67]. Nevertheless, this gradient expansion is meaningful only when the nonlocality is sufficiently weak that a truncated set of effective parameters can describe the medium response. This form of weak nonlocality underpins homogenized multipolar metasurface models [68,69].

Photonic crystals, in contrast, belong to the strong-spatial-dispersion regime [30]. In this case, nonlocality is no longer a perturbative correction to a truncated set of effective material

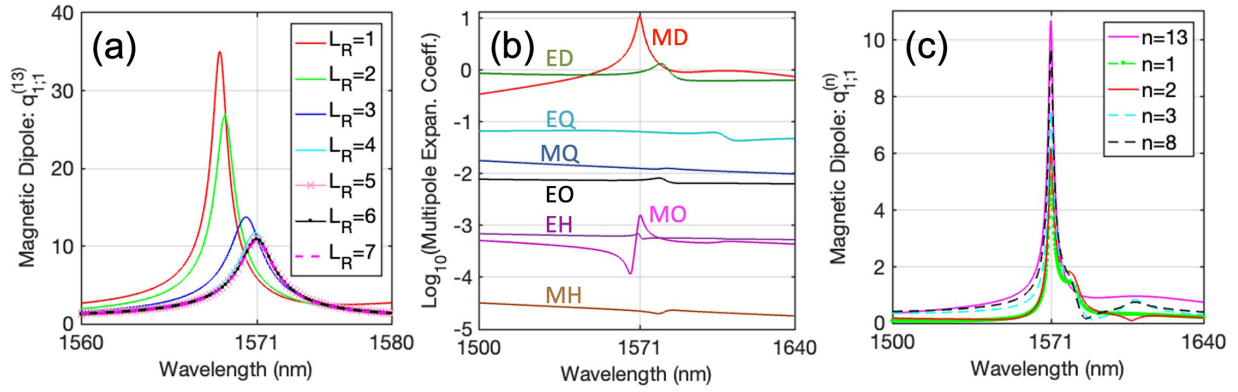


FIG. 4. Multipolar analysis of the bonding supermode B_1 excited by a z -oriented ED source placed in the gap between the 13th and 18th spheres [see Fig. 3(f)]. (a) Convergence of the MD coefficient with increasing truncation order L_R of the multipole expansion, showing that higher-order multipoles are required for convergence. (b) Multipole coefficients of the scattering field from the 13th sphere, including electric and magnetic dipole, quadrupole, octupole, and hexadecapole channels. The logarithmic scale highlights weaker higher-order contributions beyond the dominant MD response. (c) MD coefficients for the 13th, 1st, 2nd, 3rd, and 8th spheres, demonstrating that the resonance occurs collectively across the array at a common peak wavelength.

parameters, but becomes central to the governing physics. This strong nonlocality is encoded in the Bloch eigenmodes and band structure of the periodic Maxwell problem. As shown in Sec. IID, the photonic-crystal bands have the same multipolar origin as the corresponding supermodes in their finite counterparts. To further contrast this regime with the weakly nonlocal, homogenized description discussed above, we next use multipole analysis to identify how different multipolar channels contribute to the formation of nonlocal supermodes.

Before presenting a quantitative analysis, we briefly outline the multiple-scattering formulation describing long-range radiative coupling among a group of spheres. Supermodes are excited by a z -oriented ED source located at the origin O , with dipole moment $\mu_{1,0}$. The scattered field from the n th sphere, centered at O_n , is represented by electric and magnetic multipole expansion coefficients $p_{l'm'}^{(n)}$ and $q_{l'm'}^{(n)}$, respectively. The multipolar responses of the spheres are coupled through the following equations [54]:

$$p_{l'm'}^{(n)} = a_{l'}^{(n)} \left(A_{l'm'}^{1,0}(\overrightarrow{OO_n})\mu_{1,0} + \sum_{u \neq n} \sum_{l=1}^{L_u} \sum_{m=-l}^l [A_{l'm'}^{lm}(\overrightarrow{O_uO_n})p_{lm}^{(u)} + iB_{l'm'}^{lm}(\overrightarrow{O_uO_n})q_{lm}^{(u)}] \right), \quad (2)$$

$$q_{l'm'}^{(n)} = b_{l'}^{(n)} \left(-iB_{l'm'}^{1,0}(\overrightarrow{OO_n})\mu_{1,0} + \sum_{u \neq n} \sum_{l=1}^{L_u} \sum_{m=-l}^l [A_{l'm'}^{lm}(\overrightarrow{O_uO_n})q_{lm}^{(u)} - iB_{l'm'}^{lm}(\overrightarrow{O_uO_n})p_{lm}^{(u)}] \right). \quad (3)$$

Here, $A_{l'm'}^{lm}$ and $B_{l'm'}^{lm}$ are translation coefficients obtained from the addition theorem [23]. These coefficients transform outgoing multipoles from sphere u into incident multipolar fields at sphere n , thereby describing wave-mediated interactions between the induced multipolar responses of different spheres. In contrast to the weakly nonlocal homogenized description above, this formulation keeps the discrete resonator degrees of freedom explicitly. Nonlocality is therefore described through long-range radiative coupling between spatially separated spheres, including higher-order multipolar channels beyond the dipolar and quadrupolar orders.

We now turn to quantitative results for the 5×5 array of identical spheres (Fig. 4). Because all spheres have the same radius R , the same truncation order is used for all spheres, i.e., $L_u = L_R$ as defined in Eq. (1). Figure 4(a) shows the convergence of the MD response as higher-order multipoles are included in the coupled-multipole model defined by Eqs. (2) and (3). The MD spectrum associated with the central sphere

converges only when the truncation order reaches $L_R = 5$. Increasing L_R introduces additional multipolar radiation channels, which progressively broaden the resonant peak. This convergence behavior highlights a key distinction from the weak-nonlocality regime of homogenized multipolar metasurfaces, even though both descriptions use multipole expansion. In weakly nonlocal homogenized models, spatial dispersion is usually treated as a perturbative correction to an *effectively continuous response* and is often captured by low-order multipolar terms, such as dipolar and quadrupolar contributions. In the present system, in contrast, the *discrete radiative coupling* between resonators requires higher-order multipoles for convergence. The collective resonance nevertheless remains MD dominated.

We examine the multipolar composition of the resonance by plotting representative coefficients from the 13th sphere in Fig. 4(b): MD, ED, electric quadrupole (EQ), magnetic quadrupole (MQ), electric octupole (EO), magnetic octupole

(MO), electric hexadecapole (EH), and magnetic hexadecapole (MH). Each multipole order l contains $2l + 1$ modes ($m = -l, \dots, l$); here, one representative coefficient is shown for clarity. MATLAB code for calculating all multipole coefficients is publicly available [55]. Because the higher-order multipoles are off-resonant in this wavelength range, their amplitudes remain orders of magnitude smaller than those of the dominant lower-order channels.

The broad spectral overlap between the MD and ED responses (Fig. 2) is also evident in Fig. 4(b). Near 1500 nm, the ED contribution becomes comparable to, and eventually exceeds, the MD response as the system approaches the regime of collective ED resonances. In contrast, the higher-order multipoles are substantially weaker and vary smoothly across the wavelength range, without exhibiting sharp resonant features near the MD-dominated supermode. They therefore act mainly as additional radiative channels, increasing leakage and broadening the MD-dominated resonance.

For spherical resonators, supermodes associated with different Mie resonances, such as ED and MD, typically occur in distinct wavelength ranges. Each mode responds differently to geometric modifications, depending on its spatial overlap with the perturbation. By judicious symmetry breaking, the resonator geometry can be engineered to tune the spectral overlap between different Mie channels. Such tuning can lead to hybridization of nonlocal supermodes, resulting in large Q factors and enhanced Purcell effects [66].

Because supermodes arise from distributed scattering across the array, the field enhancement is strongest near the central unit cell [Fig. 3(e)]. Figure 4(c) shows the MD magnitudes for five representative spheres marked in Fig. 3(f). The amplitude decreases with distance from the center, while the resonance wavelength remains identical for all selected spheres. This behavior confirms that the supermode originates from phase-coherent multipolar coupling across the array, rather than from independent localized resonances of individual scatterers.

These results show that accurate modeling of the nonlocal supermode requires higher-order multipoles, even for sub-wavelength unit cells. Although these higher-order channels are not the dominant resonant contributions, they provide additional pathways for radiation and therefore influence the resonant linewidth. In the next section, we examine how symmetry and symmetry breaking control the formation of these supermodes and their Q factors.

III. IMPACT OF SYMMETRY BREAKING ON MIE-TRONICS SUPERMODES

Symmetry breaking of unit cells has been extensively explored for engineering nonlocal diffractive bands. Its impact on collective resonances in finite metasurfaces, however, remains far less understood. Here, we show that symmetry breaking not only modifies the coupling between Mie-tronics supermodes and external radiation but can also produce counterintuitive Q -factor enhancement—opposite to the trends predicted by infinite-lattice Bloch theory. We also show that the distinct field distributions of Mie modes provide a systematic basis for designing nonlocal metasurfaces with tailored confinement, coupling, and polarization properties.

A. Robustness of antibonding supermodes against symmetry breaking

Figure 5 summarizes the effects of horizontal and vertical symmetry breaking on Mie-tronics supermodes. Horizontal symmetry breaking is introduced by replacing square unit cells with T-shaped ones, which preserve the lattice vectors but reduce the unit-cell symmetry. Vertical symmetry breaking is implemented by introducing a quartz substrate. For the freestanding square lattice, two well-separated resonant bands appear [Fig. 5(a)], corresponding to the antibonding (A_1) and bonding (B_1) supermodes at the band edges. The magnetic-field distribution of A_1 [Fig. 5(b)] confirms its MD character, consistent with the sphere array (Fig. 3). Introducing the substrate strongly perturbs B_1 , whose field distribution extends into the surrounding environment and couples efficiently to substrate-supported channels. In contrast, A_1 remains largely unaffected because its fields are tightly confined within the unit cells, limiting interaction with the substrate.

The vortexlike far-field pattern of A_1 [Fig. 5(c)] indicates that its energy leaks primarily into oblique directions away from the surface, with weaker radiation along the in-plane directions (y and z axes). Compared with the 5×5 sphere array (Fig. 3), the 9×9 square array supports more radiation channels, reflecting stronger in-plane multiple scattering in larger systems. This enhanced scattering leads to a higher Purcell factor in Fig. 5(a). We note that the antibonding supermode A_1 corresponds to a guided resonance above the light line, as discussed in Fig. 3. In contrast, the bonding supermode B_1 behaves as a guided mode with mainly in-plane radiation. Guided resonances couple efficiently to free-space radiation and give rise to diffraction bands. In the following, we therefore focus on the antibonding supermode A_1 .

Figures 5(d)–5(f) present the response of the T-shaped array. Transforming the square unit cell into a T-shape preserves the lattice but breaks its in-plane symmetry, producing a blue shift in the Purcell spectrum [Fig. 5(d)]. Both A_1 and B_1 supermodes persist in the freestanding configuration and exhibit a substrate response similar to that of the square array. The magnetic near-field distribution [Fig. 5(e)] and far-field pattern [Fig. 5(f)] highlight a key effect of in-plane symmetry breaking: Radiation is redistributed from OCs into ICs. This redistribution increases the photon dwell time within the metasurface, enhancing the Q factor and setting the stage for the results discussed next in Fig. 6.

B. Q -factor enhancement via symmetry breaking

Figure 6(a) shows the antibonding Bloch bands above the light line for the four unit-cell geometries. The calculated band-edge wavelengths agree well with the corresponding A_1 supermodes (the hole-array case is detailed in the accompanying paper [66]). Modifying the unit-cell geometry shifts the band edges and alters the dispersion. Notably, the square and T-shaped lattices exhibit narrow spectral widths (~ 20 nm), nearly tenfold narrower than that of the hole lattice (179 nm). This result highlights the capability of geometrical design to engineer band dispersion with high precision.

Near the Γ point [Fig. 6(b)], the Bloch-mode Q factor diverges for the sphere, square, and hole lattices, reflecting a BIC at the Γ point. In contrast, the T-shaped lattice exhibits

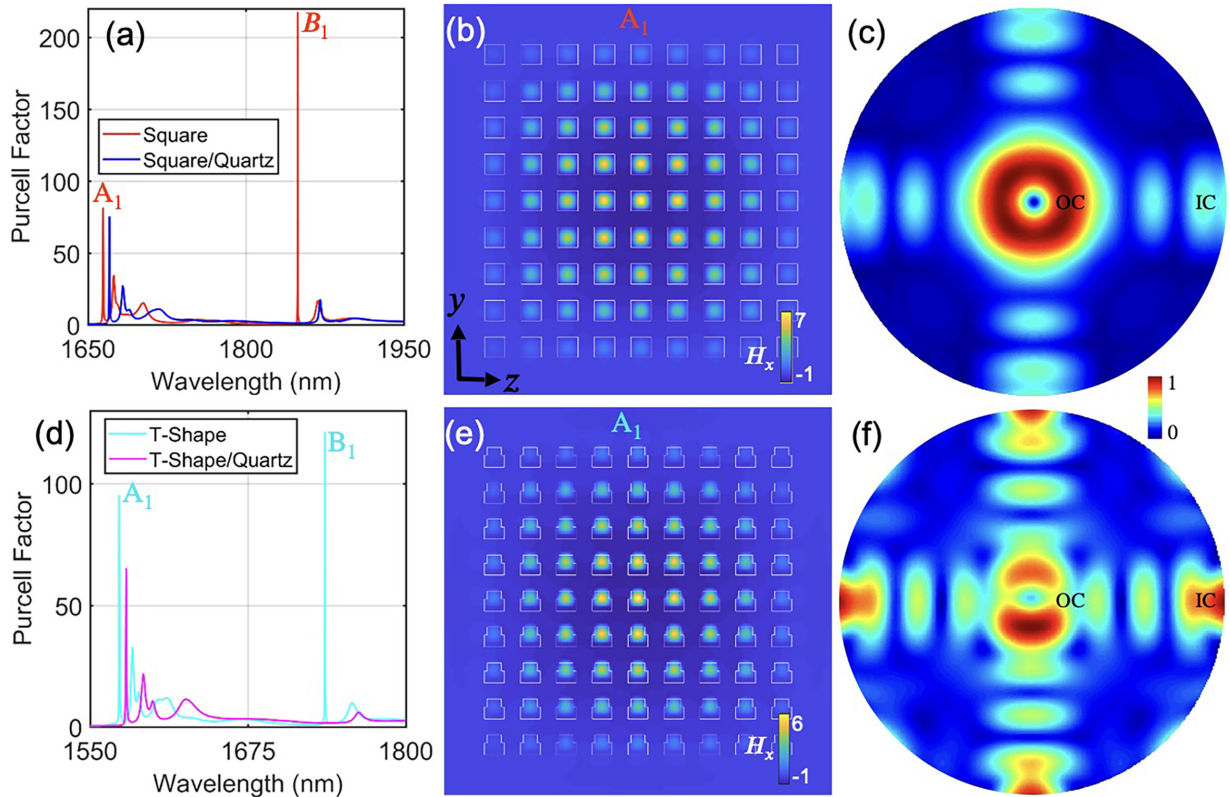


FIG. 5. Symmetry breaking enhances in-plane multiple scattering, preserves antibonding supermodes, and suppresses bonding counterparts in square and T-shaped arrays. (a) Purcell-factor spectrum for a magnetic dipole m_x at the center of a 9×9 array of square unit cells. The presence of a quartz substrate (refractive index ≈ 1.45) suppresses the bonding supermode B_1 , while the antibonding supermode A_1 remains nearly unaffected. (b), (c) Magnetic near-field distribution and vortexlike far-field pattern of A_1 . The strong oblique radiation is consistent with its guided-resonance character. (d)–(f) Same as panels (a)–(c) but for the T-shaped array. The far-field pattern highlights enhanced in-plane multiple scattering along in-plane channels (ICs) vs suppressed oblique channels (OCs).

a finite Q factor, consistent with the conventional picture in which unit-cell symmetry breaking introduces radiative leakage and reduces confinement [70].

Remarkably, the opposite trend emerges in finite metasurface arrays: Symmetry breaking enhances optical confinement by strengthening in-plane multiple scattering. As shown in

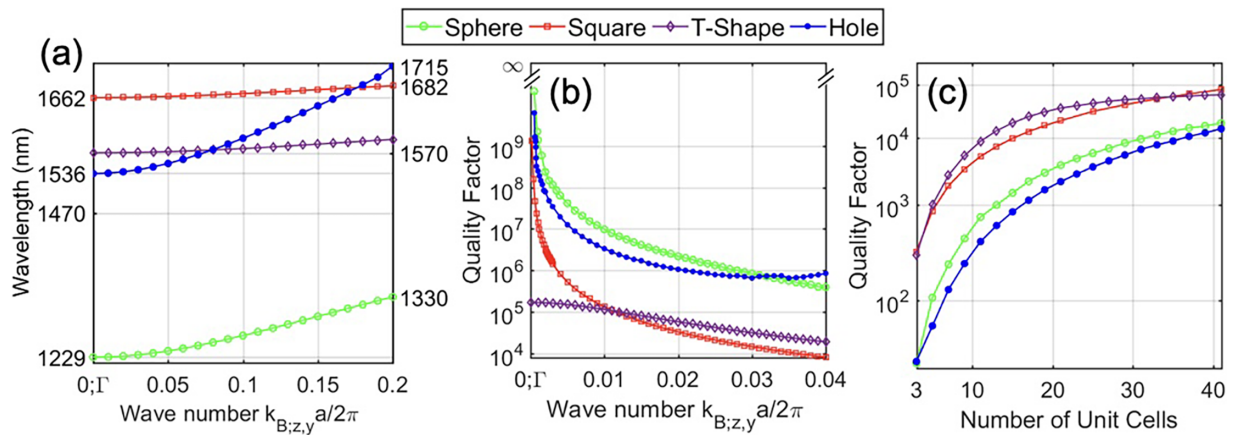


FIG. 6. Symmetry breaking lowers the quality factors of Bloch modes in infinite lattices but enhances those of supermodes in finite metasurface arrays. (a) Antibonding Bloch bands above the light line for the four unit-cell types. The band for the sphere unit cell is partially replotted from Fig. 3(d) over a reduced k -space range ($k_{B,y} = k_{B,z} \leq 0.2$) for comparison. (b) Quality factors of Bloch modes at Γ : Symmetry breaking reduces the divergent value of the square lattice to a finite one in the T-shaped lattice, while sphere and hole lattices show much higher values. (c) Quality factors of supermodes in finite arrays: Symmetry breaking boosts the quality factors of T-shaped arrays beyond that of square arrays for sizes from 5×5 to 33×33 , with both outperforming sphere and hole arrays.

Fig. 6(c), the T-shaped array achieves the highest Q factor for sizes from 5×5 to 33×33 , reversing the trend predicted by infinite-lattice Bloch analysis. This enhancement results from radiation redistribution, as discussed in the previous subsection [Figs. 5(c) and 5(f)]. For very large arrays, however, the proliferation of radiation pathways eventually saturates the Q -factor enhancement. This interplay between symmetry breaking and finite-size effects underscores the need to move beyond Bloch-wave descriptions when designing nonlocal metasurfaces for high- Q performance.

Mie-tronics therefore provides practical guidelines for designing finite nonlocal metasurfaces. This design framework is particularly relevant for the miniaturization and integration of light sources and detectors in emerging photonic platforms [71]. In the next section, we show that the modifications of unit-cell geometry and symmetry reshape the diffractive bands of these metasurfaces.

IV. DIFFRACTIVE NONLOCAL METASURFACES

The ongoing revolution in artificial intelligence has renewed interest in analog optical computing, driven by its potential for massive parallelism and ultralow energy consumption. Optical information processing, however, is not new. It has a long history in optical encryption, image processing, and early forms of optical computing [72,73]. Around the turn of this century, digital optical computing attracted significant attention. However, progress was ultimately limited by the weak optical nonlinearities of conventional materials, which hindered the realization of practical optical logic gates.

The present resurgence stems from the recognition that many artificial-intelligence tasks rely on statistical rather than logical operations, making them naturally suited to analog optical systems [74,75]. Within this context, diffractive nonlocal metasurfaces have emerged as a promising platform for high-dimensional analog computation. Their collective optical modes, extending across the entire structure, enable both deep miniaturization of on-chip architectures and performance beyond that of conventional optical components [34,76–78]. This section develops a Mie-tronics perspective on diffractive nonlocal metasurfaces, providing a physical design strategy for engineering their collective diffractive responses.

A. Mie-resonance origin of nonlocal diffractive bands

Diffractive nonlocal metasurfaces arise from the interaction between an incident plane wave and an idealized infinite lattice, as schematically illustrated in Fig. 7(a). The plane wave coherently drives all unit cells in phase, exciting their multipolar resonances at specific resonant frequencies (or wavelengths). Each unit cell therefore behaves as a nanoscale resonator supporting discrete multipolar modes, most prominently the ED and MD resonances.

For an isolated sphere, the Cartesian multipoles along the x , y , and z directions are spectrally degenerate. Coupling within an array lifts this degeneracy, splitting each vectorial multipole into several collective supermodes. For conceptual clarity, we focus on the x -oriented MD resonance throughout

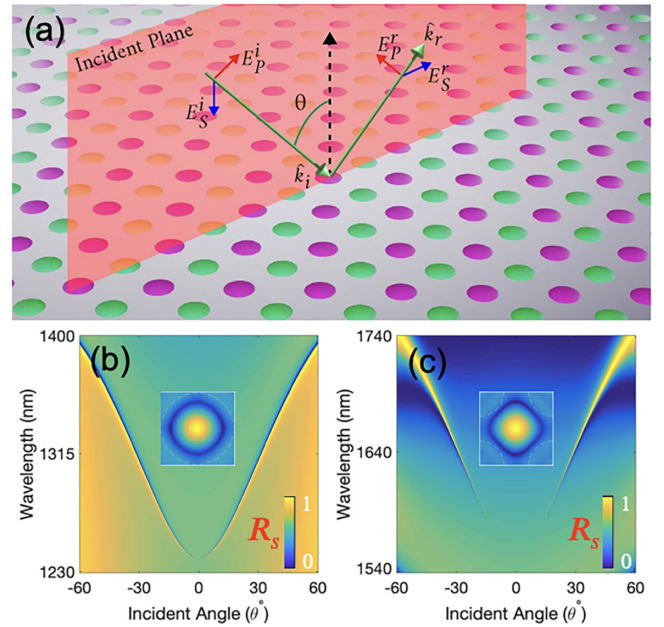


FIG. 7. Multipolar origin of diffractive nonlocal bands. (a) S- or P-polarized plane wave incident on a photonic-crystal slab. (b) Reflectance of an S-polarized plane wave from the sphere metasurface, revealing a diffractive band. The inset shows the magnetic field component $|H_x|$ at an incident angle of 5° . The band vanishes at normal incidence due to symmetry mismatch between the incident field and the eigenmode. (c) Same as panel (b) but for the hole metasurface. The inset highlights that the diffractive band originates from the MD resonance.

this work, although dipoles oriented along y or z can also form nonlocal bands under appropriate excitation.

Higher-order multipoles ($l \geq 3$) can likewise produce collective supermodes whose Bloch bands lie entirely above the light line [31]. Within diffraction theory, such higher-order bands couple to multiple radiation continua, making them particularly relevant for diffractive nonlocal metasurfaces. Their spectra, however, are densely packed and more difficult to control. Consequently, most analyses—including the present one—focus on the two lowest-order multipoles, MD and ED, whose well-separated resonances allow controlled dispersion engineering through unit-cell geometry.

In coupled arrays, both dipole and quadrupole modes split into bonding and antibonding branches. Within diffraction theory, only the antibonding supermodes couple efficiently to the radiation continua, giving rise to observable diffractive bands. Because dipole and quadrupole resonances occur at relatively low frequencies, their associated diffraction bands are spectrally well separated and tunable through geometry. Figures 7(b) and 7(c) illustrate this effect: Reshaping the spherical unit cell into a hole-type geometry transforms the MD-based metasurface from reflection-dominated to transmission-dominated operation. The near-field H_x distributions obtained from RCWA (insets) confirm that these diffractive bands originate from the underlying MD resonance.

A key feature of these bands is their disappearance at normal incidence ($\theta = 0^\circ$), commonly referred to as

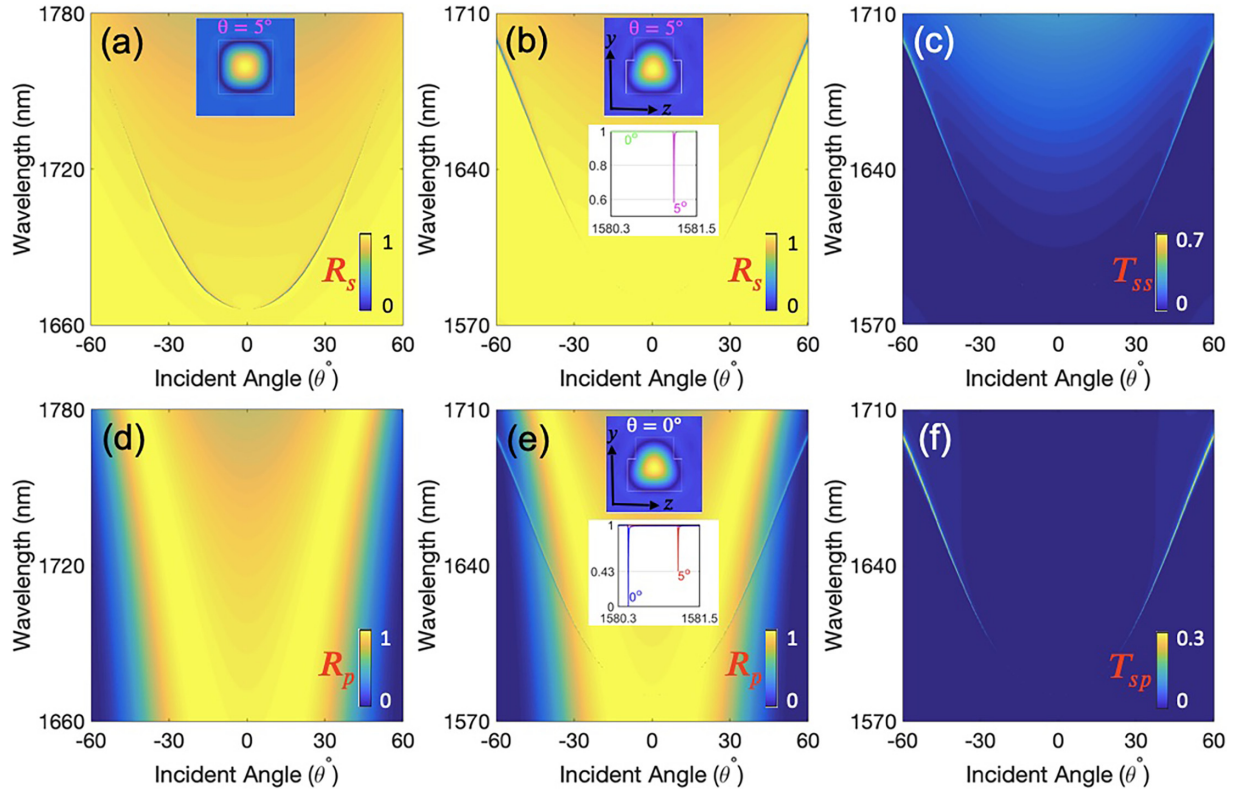


FIG. 8. Symmetry breaking enables polarization conversion in nonlocal metasurfaces. (a), (b) Diffractive nonlocal bands in square and T-shaped metasurfaces under S-polarized illumination. Insets show $|H_x|$ at an incident angle of 5° ; both bands vanish at normal incidence, as illustrated in the lower inset of panel (b). (c), (f) Transmittance spectra for S-polarized excitation of the T-shaped metasurface, revealing partial conversion from S to P polarization. (d) Reflectance of a P-polarized wave incident on the square metasurface shows no MD-based diffraction band. (e) Reflectance of a P-polarized wave incident on the T-shaped metasurface, demonstrating that symmetry breaking enables excitation of the MD-based diffraction band even at normal incidence. The upper inset shows $|H_x|$ at normal incidence ($\theta = 0^\circ$). The lower inset shows a sharp drop in reflectance from unity to zero at the eigenmode wavelength. In contrast, for oblique incidence ($\theta = 5^\circ$), the reflectance decreases from unity to ≈ 0.43 .

symmetry-protected BICs. Although such BICs are typically associated with nondegenerate diffraction bands, their physical origin is often left unexamined. Mie-tronics provides a clear interpretation by considering how electromagnetic coupling occurs through either electric or magnetic field components [79].

For the x -oriented MD mode, magnetic coupling with a P-polarized wave is forbidden because the magnetic field of P polarization lies entirely in-plane (yz plane) and is therefore orthogonal to the MD moment. Electric coupling is likewise prohibited: The in-plane electric field forms a circular displacement current that is antisymmetric with respect to the P-polarization axis (z) [Fig. 3(c)], resulting in a zero coupling coefficient. Symmetry breaking can lift this restriction, as demonstrated later for the T-shaped unit cell.

In contrast, an S-polarized wave possesses a magnetic-field component along the normal direction (x) for oblique incidence. Consequently, S polarization can excite the MD resonance except at normal incidence ($\theta = 0^\circ$), where the magnetic field has no x component. This symmetry argument explains the polarization selectivity of the MD resonance and clarifies the origin of symmetry-protected BICs in diffractive nonlocal metasurfaces.

While this BIC phenomenon is often described in terms of symmetry selection rules, its physical origin can be understood more fundamentally from the interference of multipolar currents within the unit cell. When the MD eigenmode is excited by a dipole source, multipolar displacement currents are induced in the symmetric unit cell. These currents radiate destructively along the surface normal, forming an eigenmode that cannot couple to a normally incident plane wave [31]. This interference mechanism also explains how symmetry breaking enables coupling between P-polarized light and the MD eigenmode, as demonstrated below.

B. Polarization conversion via symmetry breaking

Breaking in-plane symmetry introduces additional coupling channels that relax polarization selection rules. Figures 8(a) and 8(b) compare the S-polarized reflectance of square and T-shaped metasurfaces. The near-field distributions [insets of Figs. 8(a) and 8(b)] confirm that the diffractive bands originate from the MD resonance. The square unit cell behaves similarly to other symmetric unit cells (sphere and hole metasurfaces). Near each diffractive resonance, the S-polarized reflectance (R_s) varies rapidly from unity reflection

to complete transmission. Furthermore, the diffractive band vanishes at normal incidence due to symmetry mismatch.

Another notable feature is that the S-polarized reflectance of the T-shaped metasurface does not drop from unity to zero across the resonance [inset, Fig. 8(b)], as in symmetric unit cells. This behavior indicates the participation of an additional transmission channel in the diffraction process. Because the lattice period is subwavelength, the only accessible channel is the zeroth-order P-polarized transmission. In other words, the T-shaped metasurface partially converts S-polarized incident light into P-polarized transmission. This polarization conversion is confirmed by the spectra in Figs. 8(c) and 8(f). The conversion efficiency reaches up to 30%. It depends on the degree of symmetry breaking, which governs the electromagnetic coupling between the incident light and the eigenmode.

Interestingly, despite the symmetry breaking, the diffractive band of the T-shaped array [Fig. 8(b)] still disappears at normal incidence. Because the T-shaped lattice cannot sustain collective resonances with infinite Q factor (Fig. 6), this vanishing response cannot originate from a zero-linewidth mode. Instead, no resonance is excited at $\theta = 0^\circ$. This behavior follows from symmetry considerations. The T-shaped unit cell breaks symmetry only along the z axis and retains mirror symmetry along the y axis. Consequently, the eigenmode's electric field is antisymmetric along y and is incompatible with the symmetry of normal incidence. The eigenmode's magnetic field is oriented along x and is orthogonal to that of the incident wave. These symmetry mismatches forbid coupling between the eigenmode and the incident field.

The disappearance of the diffractive band at $\theta = 0^\circ$ is often attributed to the BIC at the Γ point in symmetric unit cells [Fig. 6(b)]. However, similar behavior has been observed experimentally in finite nonlocal metasurfaces [21]. This attribution can therefore be misleading. Finite structures cannot support true resonances at the Γ point with infinite Q factor. Such states strictly require perfect periodicity. The distinction lies in their physical origins. In finite systems, the disappearance originates from symmetry mismatch due to destructive interference. In contrast, the BIC at the Γ point arises from the singularity of wave summation in infinite lattices [31].

Figures 8(d) and 8(e) further illustrate the effects of symmetry breaking by comparing the P-polarized reflectance of square and T-shaped metasurfaces. P-polarized light cannot excite the MD-based diffraction band in the square array due to the same symmetry constraints discussed earlier. In contrast, the T-shaped array enables coupling between the MD eigenmode and the incident P-polarized wave at all angles. This coupling arises from the overlap of their electric fields along the z axis.

Under oblique incidence, the previously observed $S \rightarrow P$ conversion establishes a coupling between the polarization channels. By electromagnetic reciprocity, this coupling is symmetric. Excitation with P polarization therefore produces an S-polarized response. This is confirmed by the inset of Fig. 8(e), where the reflectance decreases from unity to a nonzero value (≈ 0.43). The behavior is consistent with that observed for S-polarized incidence at 5° [inset, Fig. 8(b)].

At normal incidence, however, the reflectance drops sharply from unity to zero at the eigenmode wavelength. This indicates the absence of polarization conversion. The MD

eigenmode remains symmetry forbidden from coupling to the S-polarized radiation channel along the surface normal, explaining the disappearance of the diffraction band in Fig. 8(b). Figure 8(e) shows that the diffractive band retains a finite linewidth across all incident angles, reflecting the finite Q factor of the T-shaped metasurface (Fig. 6).

These results show that engineering the symmetry of Mie modes provides a systematic strategy for controlling their coupling to plane waves and for optimizing nonlocal metasurfaces with tunable polarization filtering and conversion.

V. CONCLUDING REMARKS

Metasurfaces have become a cornerstone concept of modern photonics, offering miniaturized chip-scale platforms for advanced optical functionalities. Yet, pushing their performance beyond current limits requires explicit consideration of nonlocal interactions among their units. Conventional metasurface design largely relies on diffractive responses interpreted through Bloch-wave descriptions, where nonlocal lateral coupling is often treated as a complication rather than as a design principle. Bloch waves, originally developed for infinitely periodic lattices, underpin much of modern electronics by providing an intuitive description of scattering-free propagation of electrons in periodic systems. However, they cannot capture electron localization in real space. This limitation led to the introduction of Wannier functions [80], which now form the basis of electronic-structure calculations for finite systems [81].

Optical analogs of Wannier functions have been explored for modeling light localization. However, the exponential confinement characteristic of electronic Wannier functions does not naturally arise in electromagnetic systems. Instead, as shown in the accompanying work [66], multipole functions provide a more rigorous basis for describing light localization and collective resonances in metaphotonic structures. Moreover, the Q -factor enhancement arising from symmetry breaking as discussed in this work cannot be captured within Bloch-wave descriptions but is naturally explained within the emerging framework of Mie-tronics. In such three-dimensional, fully open systems, the Q factors of supermodes in T-shaped and square arrays exceed those in sphere and hole arrays, in contrast to the trends predicted by the Bloch-mode analyses. This finding highlights the importance of accounting for emitter-structure interactions within the Mie-tronics framework, particularly since realistic emitters (e.g., quantum dots) radiate into multiple channels.

The Mie-tronics design framework is built on two key elements: the Mie-mode basis and the multipolar coupling coefficients derived from the addition theorem. The Mie-mode basis describes the electromagnetic structure of individual meta-atoms, or unit cells of the metasurface, whereas the coupling coefficients capture both short- and long-range interactions among them. The importance of long-range interactions has been increasingly recognized in recent years [82–84]. Within this framework, judicious breaking of unit-cell symmetry enables selective control over how metasurface supermodes couple to external excitations, through electric-field overlap, magnetic-field overlap, or a hybrid of both. When one type of overlap dominates, the complementary

field component of the multipolar eigenmode arises from retardation-enabled displacement-current circulation within the high-index meta-atom. This selective coupling mechanism provides a clear interpretation of the Mie-resonant origin of nonlocal diffractive bands. Our results thus present Mie-tronics as a unified description of nonlocal responses in both finite and extended arrays.

Historically, the term Mie-tronics was coined in response to the growing interest in Mie resonances supported by sub-wavelength meta-atoms [45]. When such meta-atoms are arranged into strongly coupled resonant arrays, their collective response naturally enters the strong-spatial-dispersion regime [30]. Nevertheless, the underlying multipolar physics is not limited to strong nonlocality. It also underpins weakly nonlocal homogenized metasurface descriptions [68,69]. Mie-tronics therefore provides a first-principles design framework for a broad class of metasurfaces, where lateral optical communication channels become key design parameters. Many metasurfaces and metalenses are currently designed using unit-cell libraries that rely on quasilocal phase-response approximations obtained from diffractive simulations. Incorporating Mie-tronics scattering principles into such library-based design strategies may provide a fruitful direction for advancing metasurfaces and metalenses from laboratory demonstrations toward broader technological deployment [1].

Advances in high-performance computing now enable direct numerical solutions of Maxwell's equations for large ensembles of Mie scatterers. Building on these computational advances, Mie-tronics offers a systematic framework for analyzing and optimizing both the local behavior of individual meta-atoms and the nonlocal responses of large arrays. Extending this framework to nonlinear, quantum, and time-varying regimes of light-matter interaction will further expand the frontier of nonlocal photonics [85]. Such developments may ultimately enable multifunctional integrated photonic circuits that seamlessly integrate computation, control, and communication within a single platform.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [55].

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