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Surface Phonon Polariton-Quantum Dot Coupling in One-Dimensional Periodic Microstructures for Batch Quantum State Manipulation

Xinhua Zhang ^{1,2,*} , Yuchun Liu ¹, Xinyue Zhang ³, Lingchen Kong ¹, Cuihong Jin ¹, Yajuan Han ¹, Mengqing Jiang ¹, Shiying Qiao ¹ and Xinyan Gong ¹

¹ School of Photoelectric Engineering, Changzhou Institute of Technology, Changzhou 213032, China; liuyc@czust.edu.cn (Y.L.); konglc@czu.cn (L.K.); jinch@czu.cn (C.J.); hanyj@czu.cn (Y.H.); 23230311@czu.cn (M.J.); 23230319@czu.cn (S.Q.); 23230307@czu.cn (X.G.)

² Hangzhou Institute of Technology, Xidian University, Hangzhou, 311231, China

³ Chuzhou Zhimei Plasma Technology Co., Ltd., Chuzhou 239065, China; zhangxinyue@zmeitech.com

* Correspondence: zhangxh@czu.cn; Tel.: +86-1836-278-6575

Abstract

To explore the strong coupling between surface phonon polaritons (SPhPs) and quantum dots in one-dimensional periodic microstructures for quantum information processing, we establish a comprehensive theoretical model for SPhPs at air–polar dielectric interfaces. By rigorously deriving the dispersion relations, we reveal the decisive role of scale effects on bandgap formation: continuous spectra without bandgaps emerge at the nanoscale ($d \sim 10\text{--}100\text{ nm}$), whereas periodic modulation induces significant Bloch mode folding and tunable bandgaps ($0.5\text{--}5\text{ }\mu\text{m}$ width) at the microscale ($d \sim 1\text{--}10\text{ }\mu\text{m}$). Based on Fourier bandwidth limitations, we determine optimal channel widths ($L_y \geq 10\text{ }\mu\text{m}$) for maintaining low-loss modes with energy deviations below 1%. Through electromagnetic field quantization, we obtain analytical expressions for SPhP mode amplitudes and quantum dot transition rates. Calculations demonstrate that in micrometer-scale CsI structures, spontaneous emission rates can be modulated significantly: suppressed to <0.1 times the free-space values within bandgaps (excited-state lifetimes extended to $\sim 10\text{ ns}$) and enhanced 5–8 times at conduction band edges. Leveraging these characteristics, we propose a scheme for batch quantum state manipulation of $10^2\text{--}10^3$ arrayed quantum dots via selective excitation of specific Bloch modes using controlled laser frequency and angle, enabling parallel single-qubit gates with theoretical fidelity $> 99\%$. Compared with surface plasmon polariton schemes, our approach utilizes the low-loss infrared characteristics of SPhPs ($Q \sim 100\text{--}1000$, 1–2 orders higher) to reduce decoherence rates, offering a new pathway for room-temperature solid-state quantum computing and on-chip multi-node entanglement distribution.

Keywords: periodic microstructures; surface phonon polaritons; quantum dots; strong coupling; batch quantum state manipulation



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1. Introduction

Surface phonon polaritons (SPhPs) have attracted considerable attention in recent years due to their significant potential in nanophotonics, quantum technologies, and energy applications. Previous studies have demonstrated remarkable progress in various aspects, including imaging of hybridized localized and propagating SPhPs in SiC micropillar

metasurfaces [1], geometric control of SPhP properties through nanorod engineering [2], tunable hyperbolic metamaterials [3], active modulation of SPhP resonances via external excitation [4], ultraslow propagation with negative phase velocity in hyperbolic phonon polaritons [5], anisotropic van der Waals crystals [6], all-dielectric nonlinear nanoantennas for terahertz photonics [7], and giant enhancement of negative friction through resonant coupling between localized SPhPs and graphene plasmons [8]. However, the interaction between SPhPs and quantum dots in periodic nanostructures, the resulting coupling effects, and their applications in quantum information processing remain largely unexplored.

Surface phonon polaritons offer unique advantages over surface plasmon polaritons (SPPs), positioning them as highly promising candidates in the field of quantum information. First, SPhPs operate in the mid-infrared to terahertz range (1–30 THz), which naturally matches the low-energy transitions of semiconductor quantum dots (e.g., exciton transition energies of InAs/GaAs quantum dots around 10–50 meV). This avoids the energy mismatch often encountered when coupling SPPs in the visible range with solid-state quantum systems [9]. Second, polar dielectrics such as SiC and III–V semiconductors exhibit significantly lower optical losses than metals within the Reststrahlen band. As a result, SPhPs can achieve propagation lengths of tens to hundreds of micrometers and quality factors (Q) on the order of tens to hundreds, markedly outperforming SPPs [10].

Despite these advantages, existing SPhP-based quantum interfaces face critical limitations that hinder scalable quantum information processing. Flat-surface SPhP modes, while exhibiting low loss, lack frequency selectivity due to their continuous dispersion relation, rendering them unsuitable for addressing individual quantum dots in a dense array. Conversely, SPP-QD coupling schemes, though well-established in plasmonic nanocavities, suffer from high intrinsic metallic losses (ohmic damping) that limit quality factors to $Q \sim 10$ –100 and propagation lengths to $\sim 10 \mu\text{m}$, resulting in coherence times on the femtosecond to picosecond scale at room temperature. Furthermore, SPPs typically operate in the visible range, creating an energy mismatch with low-energy solid-state qubits (e.g., InAs/GaAs excitons at 10–50 meV). Our proposed 1D periodic microstructure overcomes these limitations by engineering photonic bandgaps (0.5–5 μm width) that enable frequency-selective Purcell modulation (enhancement at band edges vs. suppression within bandgaps) while preserving the low-loss characteristics of SPhPs ($Q \sim 100$ –1000, propagation lengths $> 100 \mu\text{m}$). This architecture uniquely combines the spectral tunability of photonic crystals with the long-range propagation of surface polaritons, providing a practical pathway for batch manipulation of 10^2 – 10^3 arrayed quantum dots. This implies that quantum interfaces based on SPhPs possess longer coherence times and higher information fidelity. Moreover, the frequency of SPhPs can be precisely tailored over a broad range through material selection by choosing different polar crystals with distinct phonon frequencies and structural design, such as periodic modulation. This flexibility enables the construction of reconfigurable quantum networks.

Therefore, investigating the characteristics of surface phonon polaritons in periodic nanostructures and their interaction with quantum dots presents significant value for quantum information applications. In this work, we first examine the dispersion relations of SPhPs in air–metal periodic structures, followed by field quantization to determine the corresponding mode amplitudes. By calculating the transition rates of quantum dots interacting with specific SPhP modes, we demonstrate the feasibility of batch-controlled quantum state manipulation—a critical requirement for quantum gate operations and quantum circuit construction. Our theoretical framework provides essential foundations for practical applications in quantum information processing and quantum computing.

However, existing research has yet to systematically address the following critical issues: how do the geometric parameters of periodic structures quantitatively modulate

the coupling strength between SPhPs and quantum dots? At what scale can periodic modulation induce a bandgap suitable for quantum state manipulation? How can the spatial distribution characteristics of periodic SPhP modes be exploited to achieve synchronous manipulation of multiple quantum dots? Addressing these questions is essential for the application of SPhPs in scalable quantum computing and quantum networks. In particular, one of the major bottlenecks in current quantum computing hardware is scalability—how to achieve parallel manipulation of hundreds to thousands of qubits while maintaining high-fidelity control over individual qubits. SPhP modes in periodic microstructures naturally exhibit spatial periodicity, with field intensity distributions forming standing-wave patterns consistent with the structural period, thereby providing a physical foundation for the batch addressing of arrayed quantum dots. To theoretically ground the quantum state manipulation schemes proposed in this work, we consider the interaction between a two-level quantum dot and the quantized SPhP field within the framework of cavity quantum electrodynamics. The total Hamiltonian of the coupled system is $H = H_{QD} + H_{SPhP} + H_{int}$, where $H_{QD} = \hbar\omega_0\sigma^+\sigma$ represents the quantum dot with transition frequency ω_0 (σ^+ and σ being the raising and lowering operators), $H_{SPhP} = \sum_Q \hbar\omega_Q a_Q^\dagger a_Q$ describes the quantized Bloch modes of the periodic SPhP structure with quasi-momentum Q , and $H_{int} = \sum_Q \hbar g_Q (\sigma^\dagger a_Q + \sigma a_Q^\dagger)$ represents the electric-dipole interaction under the rotating-wave approximation. The coupling strength $g_Q = \mu \cdot E_Q / \hbar$ depends on the quantum dot dipole μ moment and the quantized electric field amplitude E_Q of the SPhP mode.

In the weak coupling regime ($g_Q < \kappa/4$, where κ is the SPhP decay rate), the system exhibits the Purcell effect—spontaneous emission enhancement or suppression depending on the photonic density of states. This regime, described by the Wigner–Weisskopf theory (employed in Section 3), enables the batch manipulation schemes via frequency-selective modulation of decay rates. While the strong coupling regime (Jaynes–Cummings model, $g_Q > \kappa/4$) characterized by vacuum Rabi splitting is reserved for future investigation, the weak-coupling Purcell physics established here provides the essential theoretical basis for the room-temperature quantum state control proposed in this work.

The objective of this paper is to establish a comprehensive theoretical framework for SPhP–quantum dot coupling in periodic microstructures, elucidate the regulatory mechanisms of scale effects on coupling characteristics, and propose a scheme for batch quantum state manipulation based on SPhP bandgaps. Specifically, we will (1) rigorously derive the dispersion relation of SPhPs in one-dimensional periodic micro/nano-rod arrays and comparatively analyze mode characteristics at the nano- and microscales; (2) determine the optimal system dimensions for sustaining low-loss SPhP modes via Fourier analysis; (3) achieve the quantization of the SPhP field and develop a full quantum theory of the interaction between quantum dots and SPhPs; and (4) calculate the modulation of the spontaneous emission rate of quantum dots and propose a concrete scheme for parallel manipulation of multiple quantum dots using the frequency selectivity of the bandgap. This study provides a theoretical foundation and parametric guidance for SPhP-based solid-state quantum information processing.

This study aims to elucidate the unique properties of SPhPs in periodic micro/nanostructures and explore their potential applications in batch-controlled quantum state manipulation, thereby contributing novel theoretical insights and methodological approaches to future quantum technologies.

2. Characteristics of SPhPs in Periodic Structures

2.1. System of Micro/Nanorod Arrays Interacting with SPhPs

Before establishing the theoretical model (as depicted in Figure 1), it is necessary to clarify the basic assumptions of this paper and their scope of application. First, regarding the treatment of the damping factor: the imaginary part of the dielectric function of polar media represents the loss mechanism associated with phonon–photon coupling. In nanoscale structures, due to the significant reduction in mode volume, local field enhancement effects may lead to effective nonlinear losses, necessitating a correction to the strict loss-free assumption ($\gamma = 1$). However, this paper focuses on the qualitative influence of the scale effects (nano vs. micro) of periodic structures on the dispersion relation of SPhPs, rather than precise quantitative calculations of loss. At the micrometer scale ($d \geq 1 \mu\text{m}$), the mode volume of SPhPs is relatively large, and the field enhancement effect is relatively moderate, with experimentally measured SPhP propagation lengths reaching over $100 \mu\text{m}$. In this case, ignoring loss is a reasonable approximation. For the nanoscale ($d \leq 100 \text{ nm}$), we will qualitatively analyze the smoothing effect of loss on the bandgap in the discussion of the results, but a rigorous non-Hermitian mode analysis is reserved for future study.

Second, regarding the dielectric response of metal nanorods, this paper employs the Drude model to describe the dielectric function of the metal (gold), with a plasma frequency $\omega_p = 2.183 \times 10^{15} \text{ s}^{-1}$. It should be noted that when the nanorod width L_x decreases to below 10 nm , the effective dielectric function of the metal may deviate from the Drude model due to surface scattering and quantum confinement effects, manifesting as a blue shift of the plasma frequency and additional damping. The range considered in this paper is primarily $5\text{--}50 \text{ nm}$ (nanoscale) and $1\text{--}5 \mu\text{m}$ (microscale). At the lower limit of the nanoscale ($L_x = 5 \text{ nm}$), the correction due to size effects is approximately $10\text{--}15\%$, which does not affect the main conclusions of this paper regarding scale effects, but should be taken into account in experimental design.

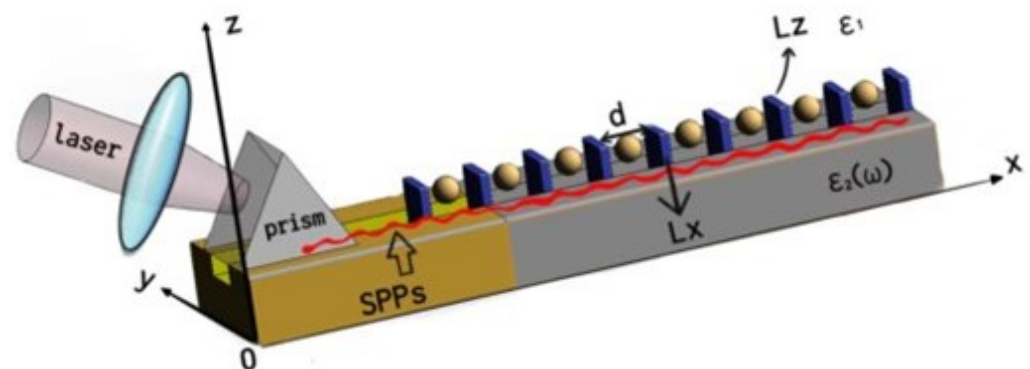


Figure 1. Model of the interaction between quantum dots and SPhP in a micro/nano periodic system: The golden spheres represent quantum dots, and the blue columns represent micro/nano rods with a transverse width of L_z and a height of L_x . SPhP is excited by the left laser beam via a prism system.

We first consider the dispersion relation of surface modes at the interface between vacuum and a polar dielectric medium (e.g., semiconductor). The polaritons are transverse magnetic (TM) modes resulting from the coupling between infrared photons and optical phonon displacements. We denote the dielectric constants of the two media as ϵ_1 and ϵ_2 , where medium 1 represents vacuum ($\epsilon_1 = 1$), and medium 2 represents a polar semiconductor with the following dielectric function [11] (It is worth emphasizing that the

dependence of the dielectric function on the wave vector direction is only applicable to isotropic materials):

$$\varepsilon_2(\omega) = \varepsilon_\infty \left(1 + \frac{\omega_L^2 - \omega_T^2}{\omega_T^2 - \omega^2 - i\omega\gamma} \right) \quad (1)$$

Here, $\varepsilon_\infty = \varepsilon(\omega)|_{\omega=\infty}$ represents the high-frequency limit of the dielectric function [12]; ω_T and ω_L are the transverse and longitudinal optical phonon frequencies, respectively, and γ is the damping factor [13]. Typically, the damping factor is small compared to SPhP frequencies (the ratios $\frac{\gamma}{\omega_T}$ or $\frac{\gamma}{\omega_L}$ range from 1% to 5% [14,15]); so, we assume $\gamma = 0$ in our calculations without significant loss of accuracy only when ω is not close to ω_T . For isotropic polar dielectrics, surface phonon polaritons are TM waves [16], meaning the magnetic field component is parallel to the interface, while the electric field $\vec{E}$ components are perpendicular to the magnetic component H_y . Due to the periodicity of the system, the wave vector component k_x may propagate in both x and $-x$ directions. Assuming that the z -component of $\vec{E}$ is a linear combination of forward and backward waves along the x -direction in the unit cell n :

$$E_{1z}^{(n)} = E_{1z} e^{-k_{z1}z - i\omega t} \left[a^{(n)} e^{ik_x(x-nd)} + b^{(n)} e^{-ik_x(x-nd)} \right] \quad (2)$$

In medium 1, the magnetic field is

$$H_{1y}^{(n)} = -\frac{\omega \varepsilon_0 \varepsilon_1 E_{1z}}{k_x} e^{-k_{z1}z - i\omega t} \left[a^{(n)} e^{ik_x(x-nd)} - b^{(n)} e^{-ik_x(x-nd)} \right] \quad (3)$$

Simultaneously,

$$E_x^{(n)} = \frac{k_{z1} E_{1z}}{ik_x} e^{-k_{z1}z - i\omega t} \left[a^{(n)} e^{ik_x(x-nd)} - b^{(n)} e^{-ik_x(x-nd)} \right] \quad (4)$$

The parameters n and d are defined as follows: n is the n -th unit cell in the periodic array, and d represents the width of each unit cell.

2.2. Dispersion Relations of SPhPs in One-Dimensional Periodic Structures

Applying standard electromagnetic boundary conditions at the interface yields the wave vector component $k_x = \frac{\omega}{c} \sqrt{\varepsilon_1 \varepsilon_2 / (\varepsilon_1 + \varepsilon_2)}$. For simplicity, we neglect the imaginary part of ε_2 by setting the damping factor $\gamma = 0$. Using Equation (1), k_x can be rewritten as

$$k_x = \frac{\omega}{c} \sqrt{\frac{\varepsilon_\infty(\omega^2 - \omega_L^2)}{(1 + \varepsilon_\infty)\omega^2 - \varepsilon_\infty\omega_L^2 - \omega_T^2}} \quad (5)$$

Thus, ω can be expressed as a function of k_x , ω_T , and ω_L [17]:

$$\omega^2(k_x) = \frac{\varepsilon_\infty}{2} \left[(1 + \varepsilon_\infty)c^2k_x^2 + \varepsilon_\infty\omega_L^2 - \Delta_{sph}^{1/2}(k_x) \right] \quad (6)$$

where

$$\Delta_{sph}(k_x) = \left[(1 + \varepsilon_\infty)c^2k_x^2 + \varepsilon_\infty\omega_L^2 \right]^2 - 1\varepsilon_\infty c^2k_x^2 (\omega_T^2 + \varepsilon_\infty\omega_L^2) \quad (7)$$

Since the right-hand side of Equation (5) must be real, and considering $\varepsilon_2 < 0$, we obtain the condition:

$$\frac{\varepsilon_\infty(\omega^2 - \omega_L^2)}{(1 + \varepsilon_\infty)\omega^2 - \varepsilon_\infty\omega_L^2 - \omega_T^2} > 0 \quad (8)$$

We find that SPhP modes occupy the frequency interval:

$$\omega_T \leq \omega(k) < \sqrt{(\varepsilon_\infty\omega_L^2 + \omega_T^2)/(1 + \varepsilon_\infty)}.$$

The frequency range given by Equation (8) can be further understood through the Lyddane–Sachs–Teller (LST) relation [18]. The LST relation indicates that the ratio of the static dielectric constant to the high-frequency dielectric constant equals the square of the ratio of the longitudinal to transverse optical phonon frequencies, namely, $\epsilon(0)/\epsilon_\infty = \omega_L^2/\omega_T^2$. This implies that the frequency range of SPhPs is entirely determined by the lattice dynamics of the material and is independent of the structural geometry. The material selection strategy shown in Table 1 is based on the following considerations: narrow-bandgap materials (such as GaSb and InSb, with $\Delta\omega/\omega_T < 7\%$, $\Delta\omega = c\Delta k_y$) are suitable for single quantum dot manipulation scenarios requiring precise frequency matching, as they can avoid crosstalk from adjacent transition energy levels; wide-bandgap materials (such as CsI and ZnS, with $\Delta\omega/\omega_T > 20\%$) are suitable for frequency-multiplexed or broadband quantum memory applications. In particular, SiC ($\Delta\omega/\omega_T \approx 19\%$) is the preferred material for experimental validation due to its excellent chemical stability, mature microfabrication processes, and the matching between its SPhP frequency ($\omega_T/2\pi \approx 24$ THz) and the typical exciton linewidth of InAs quantum dots (approximately 20–30 THz).

We now introduce the characteristic length scale $d_{02} = \frac{c}{\omega_T}$, which allows Equation (5) to be expressed in dimensionless form as

$$\tilde{k}_x = \Omega \sqrt{\epsilon_\infty \left(\Omega^2 - \frac{\omega_L^2}{\omega_T^2} \right) / \left((1 + \epsilon_\infty) \Omega^2 - \frac{\epsilon_\infty \omega_L^2}{\omega_T^2} - 1 \right)} \quad (9)$$

where $\tilde{k}_x = k_x d_{02}$, and $\Omega = \frac{\omega}{\omega_T}$. Clearly, the SPhP frequency interval depends on ω_T , ω_L and ϵ_∞ . Considering the range of $\frac{\omega}{\omega_T}$ and applying the Lyddane–Sachs–Teller (LST) relation $\omega_L^2/\omega_T^2 = \epsilon(0)/\epsilon_\infty$, we find $\frac{\omega}{\omega_T} \in \left[1, \sqrt{(1 + \epsilon(0))/(1 + \epsilon_\infty)} \right]$, where $\epsilon(0)$ is the static dielectric constant.

Table 1 presents the calculated frequency ranges for various semiconductor materials, providing guidance for selecting appropriate materials in the design of one-dimensional periodic micro/nanostructures. For example, narrow frequency bands can be achieved using GaSb or InSb, while wide frequency ranges are obtainable with CsI or ZnS.

Table 1. Frequency ranges of SPhPs in semiconductors. The values of ϵ_∞ and $\epsilon(0)$ are taken from the experimental results of Verma et al. [19] (except for the ϵ_∞ and $\epsilon(0)$ values of CsI, which are derived from Huang et al. [14]).

Material	ϵ_∞	$\epsilon(0)$	Range of $\frac{\omega}{\omega_T}$	Material	ϵ_∞	$\epsilon(0)$	Range of $\frac{\omega}{\omega_T}$
GaSb	11.11	15.69	[1, 1.01]	InP	9.61	12.55	[1, 1.13]
InSb	15.68	17.88	[1, 1.06]	ZnTe	7.28	9.67	[1, 1.11]
AlSb	10.21	12.01	[1, 1.07]	SiC	6.52	9.70	[1, 1.19]
GaAs	10.88	12.85	[1, 1.08]	ZnSe	6.20	9.20	[1, 1.19]
GaP	9.11	11.11	[1, 1.10]	CdTe	7.21	10.76	[1, 1.20]
InAs	12.25	15.16	[1, 1.10]	ZnS	5.13	8.32	[1, 1.23]
AlP	7.51	9.80	[1, 1.12]	CsI	3.0	5.33	[1, 1.26]

We now employ Bloch’s theorem to derive the second dispersion relation. According to Bloch’s theorem, the coefficients in Equations (2)–(4) satisfy the following periodicity conditions: $a^{(n+1)} = a^{(n)} e^{iQd}$, and $b^{(n+1)} = b^{(n)} e^{iQd}$, where Q denotes the one-dimensional Bloch wave vector. Applying boundary conditions at the micro/nanorods between adjacent unit cells yields the second dispersion relation:

$$\cos(Qd) - \cos(k_x d) = -\sin(k_x d) \frac{\omega_{p_{\text{nano}}}^2 L_x k_x}{2\omega^2 \epsilon_1} \quad (10)$$

Combining the first dispersion relation (Equation (5)) and the second dispersion relation (Equation (10)), we obtain

$$\begin{aligned} \cos Qd - \cos \left(\frac{\omega}{c} \sqrt{\frac{\varepsilon_{\infty}(\omega^2 - \omega_L^2)}{(1 + \varepsilon_{\infty})\omega^2 - \varepsilon_{\infty}\omega_L^2 - \omega_T^2}} d \right) = \\ - \sin \left(\frac{\omega}{c} \sqrt{\frac{\varepsilon_{\infty}(\omega^2 - \omega_L^2)}{(1 + \varepsilon_{\infty})\omega^2 - \varepsilon_{\infty}\omega_L^2 - \omega_T^2}} d \right) \frac{\omega_{p,\text{nano}}^2 L_x \sqrt{\frac{\varepsilon_{\infty}(\omega^2 - \omega_L^2)}{(1 + \varepsilon_{\infty})\omega^2 - \varepsilon_{\infty}\omega_L^2 - \omega_T^2}}}{2\omega c} \end{aligned} \quad (11)$$

From Equation (11), we can establish the relationship between Q and ω :

$$\begin{aligned} Qd = \arccos \left[\cos \left(\frac{\omega}{c} \sqrt{\frac{\varepsilon_{\infty}(\omega^2 - \omega_L^2)}{(1 + \varepsilon_{\infty})\omega^2 - \varepsilon_{\infty}\omega_L^2 - \omega_T^2}} d \right) \right. \\ \left. - \sin \left(\frac{\omega}{c} \sqrt{\frac{\varepsilon_{\infty}(\omega^2 - \omega_L^2)}{(1 + \varepsilon_{\infty})\omega^2 - \varepsilon_{\infty}\omega_L^2 - \omega_T^2}} d \right) \times \frac{\omega_{p,\text{nano}}^2 L_x \sqrt{\frac{\varepsilon_{\infty}(\omega^2 - \omega_L^2)}{(1 + \varepsilon_{\infty})\omega^2 - \varepsilon_{\infty}\omega_L^2 - \omega_T^2}}}{2\omega c} \right] \end{aligned} \quad (12)$$

The physical content of the dispersion relation (Equation (12)) reveals the dual nature of SPhP propagation in periodic structures. The term $\cos(k_x d)$ represents the phase accumulation of the surface wave across a single unit cell, while the right-hand side embodies the scattering contribution from the metallic nanorods. When the Bragg condition $Qd = m\pi$ is approached, the periodic modulation induces constructive interference of scattered waves, leading to the formation of standing wave patterns. This manifests as a divergence in the density of states at the zone boundary, which will later prove crucial for enhancing the quantum dot transition rate. The characteristic length scale $d_{o2} = \frac{c}{\omega_T}$ (Equation (9)) determines the transition from the “effective medium” regime ($d \ll d_{o2}$) to the “photonic crystal” regime ($d \sim d_{o2}$): only when the period becomes comparable to the SPhP wavelength does the zone-folding effect become pronounced enough to open measurable bandgaps.

Since we have defined the characteristic frequency $\Omega = \frac{\omega}{\omega_T}$, Equation (11) becomes

$$\cos Qd - \cos \left[\Omega S(\Omega) \frac{\omega_T d}{c} \right] = -\sin \left[\Omega S(\Omega) \frac{\omega_T d}{c} \right] \frac{\omega_{p,\text{nano}}^2}{2\Omega c \omega_T} S(\Omega) \quad (13)$$

where $S(\Omega) = \sqrt{\frac{\varepsilon_{\infty}(\Omega^2 - \omega_L^2/\omega_T^2)}{(1 + \varepsilon_{\infty})\Omega^2 - \varepsilon_{\infty}\omega_L^2/\omega_T^2 - 1}} = \tilde{k}_x/\Omega$.

From Equation (12), we can calculate the allowed range of Qd values. When $\Omega = 1$, the minimum value of Qd is

$$Qd|_{\text{start}} = \arccos \left[\cos \left[S(1) \frac{\omega_T d}{c} \right] - \sin \left[S(1) \frac{\omega_T d}{c} \right] \frac{\omega_{p,\text{nano}}^2}{2c\omega_T} S(1) L_x \right] \quad (14)$$

Similarly, for the other extreme, we obtain

$$Qd|_{\text{end}} = \arccos \left[\cos \left[\Omega_{\text{lim}} S(\Omega_{\text{lim}}) \frac{\omega_T d}{c} \right] - \sin \left[\Omega_{\text{lim}} S(\Omega_{\text{lim}}) \frac{\omega_T d}{c} \right] \frac{\omega_{p,\text{nano}}^2}{2\Omega_{\text{lim}} c \omega_T} S(\Omega_{\text{lim}}) L_x \right] \quad (15)$$

where $\Omega_{\text{lim}} = \sqrt{(\varepsilon_{\infty}\omega_L^2/\omega_T^2 + 1)/(1 + \varepsilon_{\infty})}$. Table 2 shows the values of $Qd|_{\text{start}}$ and $Qd|_{\text{end}}$ for different materials.

Table 2. Values of ω_T and ω_L are sourced from Kittel [12], Amirtharaj [20], Huang [11], et al. The units for ω_T and ω_L are 10^{13} s^{-1} .

Material	$Qd _{\text{start}}^{d=100 \text{ nm}}$	$Qd _{\text{end}}^{d=100 \text{ nm}}$	ω_T	ω_L
GaSb	0.231	0.212	1.3	1.6
InSb	0.231	1.391	3.5	3.7
GaAs	0.231	1.629	5.1	5.5
GaP	0.232	2.101	6.9	7.6
InAs	0.231	2.373	1.1	1.5
InP	0.231	2.527	5.7	6.5
SiC	0.251	2.325	11.9	17.9
ZnS	0.250	0.111	5.16	6.63
CsI	0.271	2.515	7.51	10.05

2.3. Results for Nanoscale Systems

To investigate the validity of the derived dispersion relations at the nanoscale, we select specific material parameters ω_T , ω_L and $\omega_{p_{\text{nano}}}$, assuming that the unit cell width d , nanorod width L_x , and channel width L_y are all on the nanometer scale. The $\omega - Q$ dispersion relation is then obtained from Equation (10). For example, we select CsI as the substrate material with parameters [11] $\omega_T = 7.51 \times 10^{13} \text{ s}^{-1}$ and assume $\varepsilon_1 = 1$. For the nanorods, we select gold with parameters $\omega_{p_{\text{nano}}} = 2.183 \times 10^{15} \text{ s}^{-1}$ following the Drude model.

To systematically analyze these dispersion characteristics, we present numerical results for various parameter configurations. First, varying the unit cell width d while keeping other parameters constant produces the results shown in Figure 2. As d increases, the starting point of Qd on the x -axis gradually shifts toward larger values.

Second, varying L_x while maintaining other parameters constant yields the results shown in Figure 3. This indicates that changing L_x also modifies the minimum allowable Qd value: smaller L_x corresponds to smaller Qd values. The overall trend of the curves is similar to Figure 2. (It should be emphasized that to demonstrate how the dispersion relation varies with L_x , the values of L_x do not always satisfy $L_x \ll d$.)

Furthermore, in Figure 4, we plot the dispersion relation $\frac{\omega}{\omega_T}$ versus $\frac{ck_x}{\omega_T}$ under two extreme conditions: Figure 4a, where d is large compared to the nanorod thickness, and Figure 4b, where $L_x = 0$. It is evident that in both cases with large d and $L_x = 0$, the results are essentially identical. Additionally, since there are no nanorods at the interface in these extreme cases, the results are expected—the dispersion curves of $\frac{\omega}{\omega_T}$ versus $\frac{ck_x}{\omega_T}$ resemble those of a simple vacuum–polar dielectric interface.

In summary, SPhPs in one-dimensional periodic nanostructures exhibit the following characteristics: (i) No significant frequency bandgaps are observed in the $\frac{\omega}{\omega_T} \sim Qd$ dispersion relations. Due to the fact that SPhP frequencies are typically two orders of magnitude lower than those of surface plasmon polaritons (SPPs), bandgap effects become indiscernible at the extremities of the Qd axis in nanoscale systems. (ii) The entire interval of $\frac{\omega}{\omega_T}$ lies within $\left[1, \sqrt{\left(\frac{\varepsilon_\infty \omega_L^2}{\omega_T^2} + 1\right) / (1 + \varepsilon_\infty)}\right]$. (iii) The unit cell width d determines the initial position of the dispersion curve, with an increasing d resulting in larger minimum allowable values of Qd . The nanorod width L_x exhibits similar effects on the dispersion characteristics. (iv) In the limiting cases when d is very large compared to the width L_x of the nanorod or $L_x = 0$, the dispersion relations converge to those of a bare vacuum–dielectric interface without nanorod arrays.

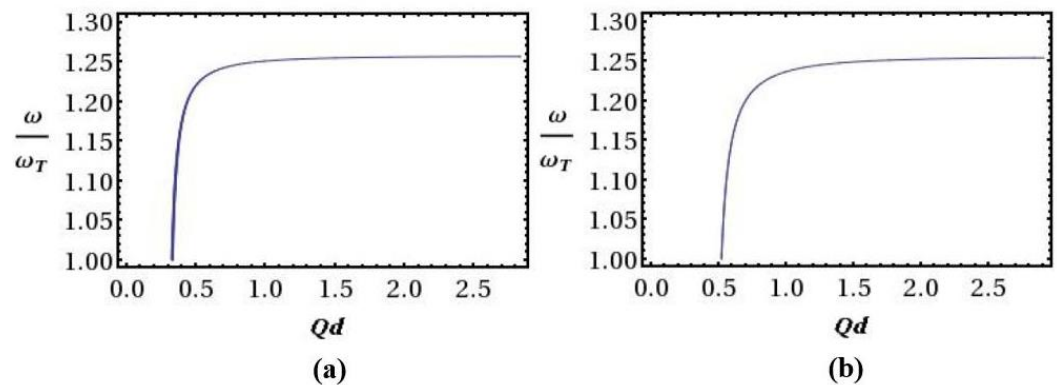


Figure 2. Plots of $\frac{\omega}{\omega_T}$ versus Qd for different cell widths. In panel (a), $d = 200$ nm; in panel (b), $d = 500$ nm. For panel (a), $Qd|_{start} = 0.338$, $Qd|_{end} = 2.836$; for panel (b), $Qd|_{start} = 0.513$, $Qd|_{end} = 2.877$.

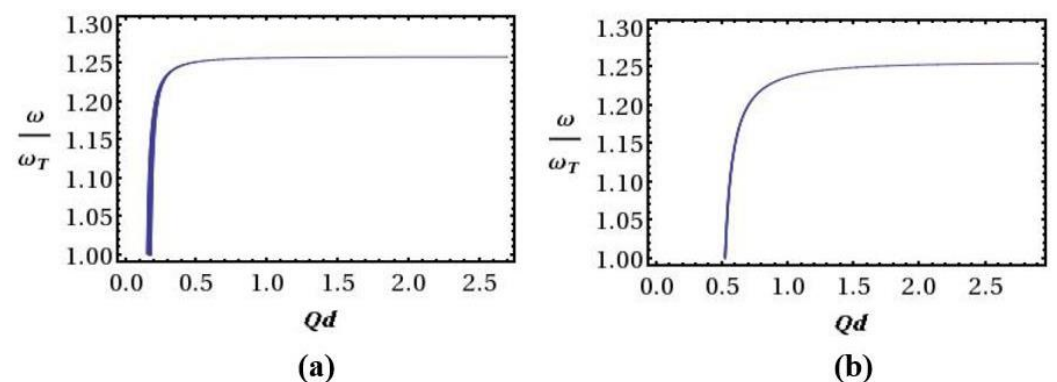


Figure 3. Plots of $\frac{\omega}{\omega_T}$ versus Qd for different cell widths. In panel (a), $L_x = 5$ nm; in panel (b), $L_x = 50$ nm. For panel (a), $Qd|_{start} = 0.193$, $Qd|_{end} = 2.780$; for panel (b), $Qd|_{start} = 0.510$, $Qd|_{end} = 2.821$.

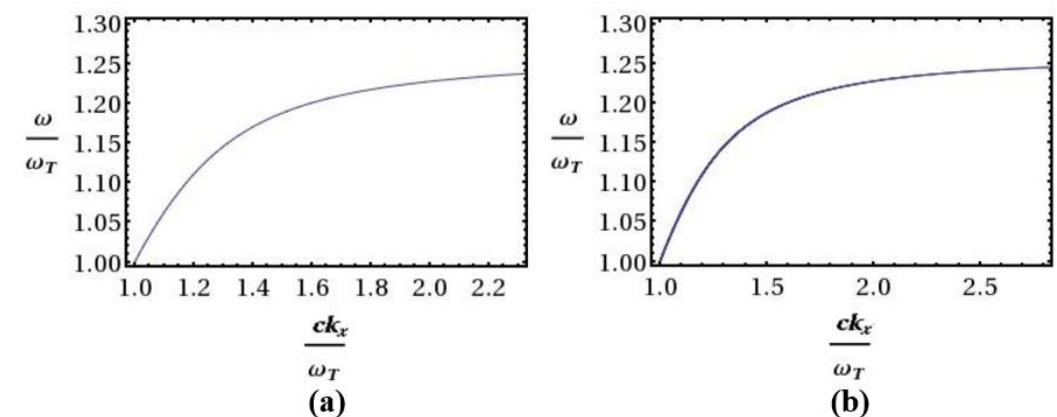


Figure 4. Dispersion curves of $\frac{\omega}{\omega_T}$ versus $\frac{ck_x}{\omega_T}$ under two extreme conditions. In panel (a), $d = 1000$ nm, $L_x = 10$ nm; in panel (b), $d = 100$ nm, $L_x = 0$ nm.

It is particularly important to note that the continuous dispersion characteristics (absence of a bandgap) of SPhPs at the nanoscale do not render them valueless for quantum information applications. On the contrary, a continuous spectrum implies a broadband response, which can be utilized to achieve broad-spectrum coupling between quantum dots and SPhPs, making it suitable for scenarios requiring rapid spontaneous emission (enhanced by the Purcell effect). However, due to the lack of a bandgap, it is impossible to leverage frequency selectivity for the individual addressing of multiple quantum dots, nor is it feasible to construct quantum gates based on photon blockade. Furthermore, at the

nanoscale, the wavevector of SPhPs can significantly exceed that of free-space light. This momentum mismatch results in extremely low far-field excitation efficiency, necessitating reliance on near-field optical techniques such as scanning near-field optical microscopy (SNOM). This increases the experimental complexity and limits the integration density. Therefore, although nanostructures offer advantages in mode compression and field enhancement, micro-scale periodic structures demonstrate superior overall performance for quantum information processing that requires batch manipulation and on-chip integration.

2.4. Results for Micrometer-Scale Systems

However, we note that the characteristic length scale $d_{O2} = \frac{c}{\omega_T}$ is on the order of 10^{-5} m. Therefore, it is essential to examine the dispersion characteristics in the 10^{-5} m regime, where system dimensions approach this characteristic length scale. Figures 5 and 6 clearly demonstrate that when d and L_x become comparable to d_{O2} , both dispersion relations, $\frac{\omega}{\omega_T}$ versus Qd and $\frac{\omega}{\omega_T}$ versus $\frac{ck_x}{\omega_T}$, exhibit the expected zone-folding phenomena characteristic of periodic structures, most notably the emergence of frequency bandgaps. By adjusting parameters such as d or L_x within appropriate ranges, we can observe how the dispersion curves evolve. Figures 6 and 7 provide examples illustrating these variations. Compared with the nanoscale regime, the micrometer-scale system exhibits distinct characteristics (Figures 5 and 6): Figure 6a,b show that a smaller L_x corresponds to smaller bandgaps, and vice versa. As expected, when L_x is sufficiently small, the bandgap disappears, which is precisely the phenomenon observed in nanoscale SPhP–quantum dot interactions. From this perspective, adjusting the microrod width enables the selection of specific frequency bands with desired widths for optimized SPhP–quantum dot interactions.

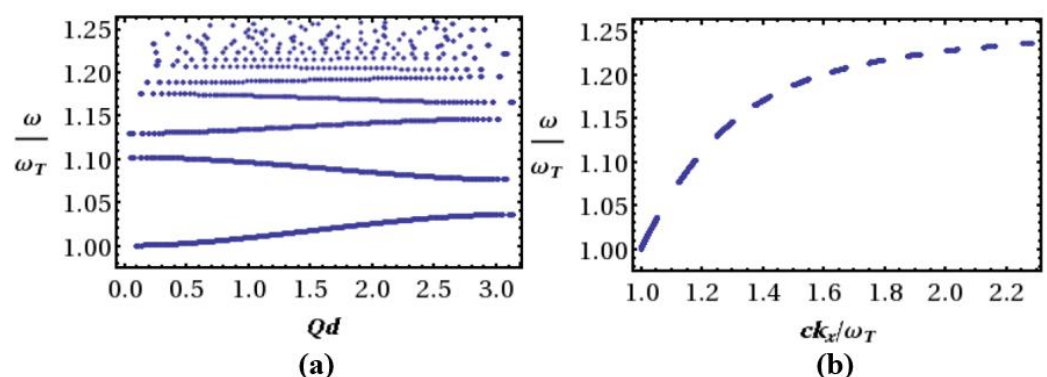


Figure 5. (a) Dispersion relation of $\frac{\omega}{\omega_T}$ versus Qd . The substrate material is CsI, with $\omega_T = \omega_{p_{\text{nano}}}$. The parameters used in the figure are $d = 10 \times 10^{-5}$ m, $L_x = 1.0 \times 10^{-5}$ m, $\epsilon_1 = 1$, indicating that the quasi-nanorods are made of doped semiconductors. As indicated by the text at the top, more and more folded branches appear as the lines approach the $\frac{\omega}{\omega_T}$ limit (The dotted lines were produced by solving the equations from a Fortran program and were plotted by Mathematica; since the roots were discrete, the plotted lines appeared as dotted shapes). (b) Schematic dispersion relation of $\frac{\omega}{\omega_T}$ vs. $\frac{ck_x}{\omega_T}$. Parameters chosen: $d = 10 \times 10^{-5}$ m, $L_x = 1.0 \times 10^{-5}$ m, $\omega_T = \omega_{p_{\text{nano}}}$.

The emergence of bandgaps in the microscale regime ($d \sim 1\text{--}10\ \mu\text{m}$) can be understood through the interplay of Bragg reflection and mode confinement. Physically, the periodic array of microrods creates a spatially modulated dielectric environment that reflects SPhP modes when the Bloch wavelength $\lambda_B = 2\pi/Q$ satisfies the Bragg condition $\lambda_B \approx 2d/m$. This forbidden frequency range (bandgap) arises from the destructive interference of forward and backward propagating Bloch waves, effectively localizing photonic states and reducing the group velocity $v_g = d\omega/dQ$ to near zero at the band edges (Figure 6).

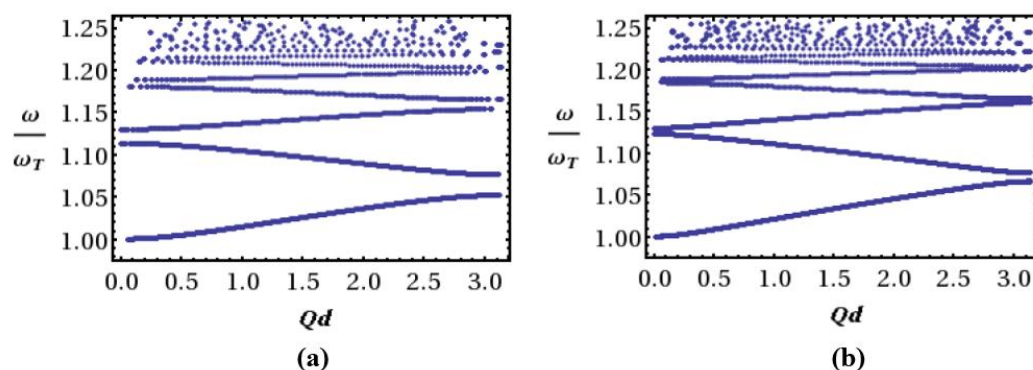


Figure 6. Schematic diagram of $\frac{\omega}{\omega_T}$ versus Qd for micro-rods of different widths. Subfigures (a,b) adopt the same parameters as Figure 5: for subfigure (a), $L_x = 5 \mu\text{m}$; for subfigure (b), $L_x = 2 \mu\text{m}$.

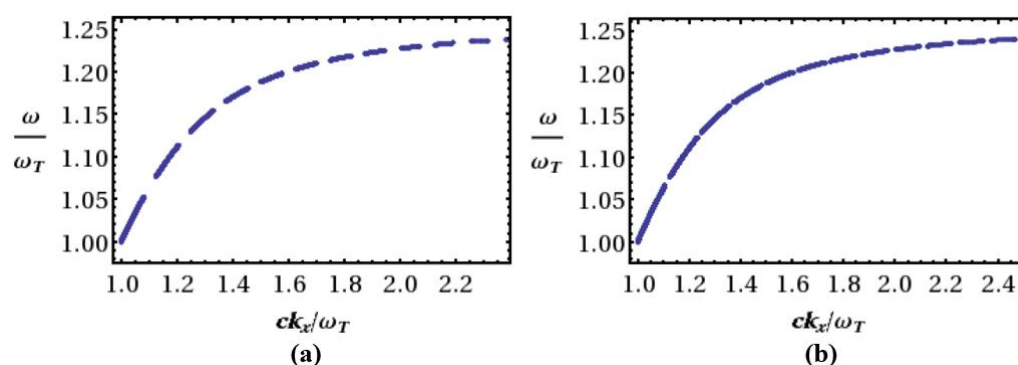


Figure 7. Plots of the relationship of $\frac{\omega}{\omega_T}$ versus $\frac{ck_x}{\omega_T}$ for micro-rods of different widths. Subfigures (a,b) adopt the same parameters as Figure 5: for subfigure (a), $L_x = 5 \mu\text{m}$; for subfigure (b), $L_x = 2 \mu\text{m}$.

The geometric parameters govern this process through distinct mechanisms:

- Periodicity d : Determines the Bragg wavelength and the number of folded Brillouin zones. Increasing reduces the spacing between adjacent bands (Equations (13)–(15)), compressing the available phase space for propagation and increasing the density of states at the band edges (van Hove singularities).
- Rod width L_x : Controls the scattering strength of the periodic perturbation. As shown in Figure 6, L_x reducing from $5 \mu\text{m}$ to $2 \mu\text{m}$ weakens the scattering potential, narrowing the bandgap width from $\Delta\omega/\omega_T \approx 8\%$ to 3% . This linear dependence arises because the perturbation Hamiltonian is proportional to the rod polarizability, which scales with L_x in the dilute limit ($L_x \ll d$).

Crucially, the group velocity dispersion $\partial\omega/\partial Q$ vanishes at the band edges (Figure 6), creating “slow-light” conditions where the SPhP mode density of states $\rho(\omega) \propto 1/v_g$ diverges. This enhancement directly amplifies the light–matter coupling rate, as the Purcell factor scales as $F_p \propto Q/V_{\text{mode}} \propto 1/(v_g \cdot A_{\text{eff}})$, where A_{eff} is the effective mode area in the transverse direction.

The frequency gaps in Figure 7a,b correspond to those in Figure 6a,b, respectively. Unlike the continuous dispersion relations in nanoscale systems (Figures 3 and 4), the micrometer-scale results (Figures 6 and 7) show segmented dispersion curves with well-defined bandgaps. As frequency increases, the frequency gaps become smaller. Comparing Figures 6 and 7, we observe that the frequency range remains identical to that in nanoscale systems: $\left[1, \sqrt{\frac{\epsilon_\infty \omega_L^2}{\omega_T^2} + 1} / (1 + \epsilon_\infty)\right]$, indicating that the frequency range is independent of system scale. It is noteworthy that as $\frac{\omega}{\omega_T}$ increases, the gradient of the dispersion

curve decreases, approaching zero, which implies slower surface mode propagation at higher frequencies.

The emergence of a bandgap for SPhPs at the microscale provides a critical tool for quantum state manipulation. Specifically, when the transition frequency of a quantum dot falls within the SPhP bandgap, the photonic density of states (DOS) is suppressed, leading to a significant inhibition of the quantum dot's spontaneous emission (with suppression factors potentially reaching $10^{-2} \sim 10^{-3}$). This prolongs the excited state lifetime to the microsecond level, creating conditions suitable for quantum memory and high-fidelity single-photon sources. Conversely, when ω_0 is positioned near the edge of the conduction or valence band, the van Hove singularity in the DOS results in enhanced spontaneous emission (with Purcell factors potentially reaching $10^2 \sim 10^3$), which can be exploited for implementing fast quantum gate operations and efficient single-photon extraction. More importantly, by adjusting the microrod width L_x , the band-edge frequency can be continuously tuned within a certain range (as shown in Figure 6, a reduction from $5 \mu\text{m}$ to $2 \mu\text{m}$ shifts the bandgap center frequency by approximately 5%). This implies that dynamic tuning of L_x (e.g., using MEMS actuation) through electro-optic or thermo-optic effects is possible, enabling real-time switching of the quantum dot–SPhP coupling strength—a fundamental capability for building reconfigurable quantum circuits.

2.5. Determination of Optimal System Dimensions via Fourier Bandwidth Constraints

As previously discussed, Fourier bandwidth limitations provide a criterion for estimating the optimal channel width required to sustain SPhP modes. Given that SPhP energies typically range from 27 to 59 meV [19,20], the uncertainty principle yields

$$\Delta E = \hbar \Delta \omega \geq \frac{\hbar c}{2 \Delta L_y} \quad (16)$$

Assuming the minimum uncertainty product, we obtain an optimal channel width of $\Delta L_y \in [1.6, 3.5] \mu\text{m}$, which represents the minimum dimension required to sustain typical SPhP modes with negligible energy uncertainty. Furthermore, this allows selection of appropriate channel widths where energy deviations can be neglected; for example, when $\Delta L_y = 10 \mu\text{m}$, $\Delta E \approx 1 \text{ meV}$, corresponding to energy deviations less than 1%. Here, we establish that our theoretical analysis is valid exclusively for micrometer-scale systems. The above analysis provides clear guidelines for the design of experimental systems. First, the choice of channel width L_y requires a trade-off between mode confinement and energy precision: with $L_y = 10 \mu\text{m}$, the energy resolution $\Delta E \approx 1 \text{ meV}$ is sufficient to resolve adjacent Bloch modes (typical mode spacing is approximately 5–10 meV), while ensuring effective transverse confinement of the SPhPs. Second, the longitudinal length L_z (along the y -direction) should be much larger than the propagation length of SPhPs (approximately $100 \mu\text{m}$) to ensure that the quasi-infinite length approximation for the mode in the longitudinal direction holds, avoiding mode hybridization caused by end-face reflections. Finally, the number of array periods N should satisfy $N \times d \gg 2\pi/Q$ to guarantee the applicability of Bloch's theorem. For typical parameters ($d = 10 \mu\text{m}$, $Q \approx 0.2\pi/d$), the condition $N \geq 10$ is sufficient, which is readily achievable in microfabrication processes.

3. Quantization of SPhPs and Transition Rates of Quantum Dots

3.1. Quantization Results of SPhPs

To quantize the SPhP modes, we perform a detailed derivation and obtain the amplitude of the electric field component E_{1z} as

$$|E_{2z}|^2 \left(a^{(0)^2} + b^{(0)^2} \right) = \frac{\hbar k_{z1} k_x^2 c^2}{2L_y \omega \epsilon_0 \epsilon_1^2} \frac{1}{\left[3 + \epsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2 (\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} \quad (17)$$

Furthermore, we obtain

$$E_{1z}^2 a^{(0)^2} = \frac{\hbar k_{z1} k_x^2 c^2}{2L_y \omega \epsilon_0 \epsilon_1^2} \frac{A^2}{A^2 + B^2} \frac{1}{\left[3 + \epsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2 (\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} \quad (18)$$

and

$$E_{1z}^2 b^{(0)^2} = \frac{\hbar k_{z1} k_x^2 c^2}{2L_y \omega \epsilon_0 \epsilon_1^2} \frac{B^2}{A^2 + B^2} \frac{1}{\left[3 + \epsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2 (\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} \quad (19)$$

where $A = 1 - e^{iQd+ik_x d}$, and $B = 1 - e^{iQd-ik_x d}$. Through this quantization procedure, we have determined the amplitudes of the surface modes, enabling calculation of the transition rates when SPhPs interact with two-level quantum dots. The quantization of SPhPs follows the standard procedure for electromagnetic field quantization. However, due to the presence of the periodic structure, special attention must be paid to the orthonormalization of the Bloch modes. In the z-direction (perpendicular to the interface), the SPhP field decays exponentially and lacks a discrete mode spectrum; therefore, an effective mode length $L_{eff} = 1/2k_{z1}$ must be introduced for the normalization volume. In the x-direction (the periodic direction), the orthogonality of the Bloch modes is guaranteed by the integral $\int_{cell} u_Q^*(x) u_{Q'}(x) dx = \delta_{QQ'}$, where $u_{Q(x)}$ is the periodic part of the Bloch function. Since we are considering a finite array (with N periods), the Bloch wavevector Q is discretized as $2\pi m/(Nd)$, with $m = 1, 2, \dots, N$. In the thermodynamic limit ($N \rightarrow \infty$), the discrete spectrum becomes continuous. In this case, the summation in the transition rate calculation must be replaced by an integral: $\sum_Q \rightarrow \frac{Nd}{2\pi} \int dQ$. This is precisely the origin of the integral form in Equation (20). It is important to emphasize that this treatment assumes the weak coupling regime between the quantum dot and the SPhPs (i.e., the spontaneous emission rate is much smaller than the linewidth of the SPhP mode). Under this condition, the Wigner–Weisskopf approximation is valid, and the transition rate formula (Equation (24)) holds. For the strong coupling regime (such as vacuum Rabi splitting in microcavities), the Jaynes–Cummings model must be employed, which is reserved for future study.

3.2. Transition Rates for SPhPs Interacting with Quantum Dots

Here, we transform the frequency delta function using the density of states:

$$\begin{aligned} \Gamma &= \int \frac{2\pi}{\hbar^2} |\langle i | -\vec{d} | f \rangle \cdot \vec{E}_e(Q)|^2 \delta(\omega_0 - \omega) dQ \\ &= \int \frac{2\pi}{\hbar^2} |\langle i | d_x | f \rangle \cdot E_{1x} + \langle i | d_z | f \rangle \cdot E_{1z}|^2 \delta(\omega_0 - \omega) dQ \end{aligned} \quad (20)$$

Here, $\delta(\omega - \omega_0) = \frac{\delta(Q - Q_0)}{\left| \frac{d\omega}{dQ} \right|} = \left| \frac{dQ}{d\omega} \right| \delta(Q - Q_0)$. Consequently, Equation (20) becomes

$$\Gamma = \int \frac{2\pi}{\hbar^2} |\langle i | d_x | f \rangle \cdot E_{1x} + \langle i | d_z | f \rangle \cdot E_{1z}|^2 \left| \frac{dQ}{d\omega} \right| \delta(Q - Q_0) dQ \quad (21)$$

From Equation (10), we obtain

$$\frac{dQ}{d\omega} = \frac{d\sin(k_x d)k'_x + \frac{\omega_{\text{p nano}}^2 L_x}{2\varepsilon_1 \omega^3} \left[d\omega k_x \cos(k_x d)k'_x + \omega \sin(k_x d)k'_x - 2\sin(k_x d)k_x \right]}{\sin(Qd)d} \quad (22)$$

where

$$k'_x = \frac{1}{c} \left[\sqrt{\frac{\varepsilon_\infty(\omega^2 - \omega_L^2)}{(1 + \varepsilon_\infty)\omega^2 - \varepsilon_\infty\omega_L^2 - \omega_T^2}} + \frac{\omega^2 \varepsilon_\infty (\omega_L^2 - \omega_T^2)}{\sqrt{\varepsilon_\infty(\omega^2 - \omega_L^2) [(1 + \varepsilon_\infty)\omega^2 - \varepsilon_\infty\omega_L^2 - \omega_T^2]^3}} \right] \quad (23)$$

Thus, the transition rate takes the form:

$$\Gamma = \frac{2\pi}{\hbar^2} \left| \langle i | d_x | f \rangle \cdot E_{1x} + \langle i | d_z | f \rangle \cdot E_{1z} \right|^2 \left| \frac{dQ}{d\omega} \right|_{Q=Q_0} \quad (24)$$

With the assistance of Equations (2) and (4), we now obtain results for the transition rates of two-level atoms with specific initial and final states. When the dipole is parallel to the interface surface, the transition rate is

$$\Gamma_{\parallel} = |\langle i | d_x | f \rangle|^2 \frac{2c^2 k_{z1}^3 e^{-2k_{z1}z}}{\hbar L_y \varepsilon_0 \varepsilon_1^2 \omega} \frac{1}{\left[3 + \varepsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2(\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} f(A - B) \left| \frac{dQ}{d\omega} \right|_{Q=Q_0} \quad (25)$$

When the dipole is perpendicular to the interface surface, we have

$$\Gamma_{\perp} = |\langle i | d_z | f \rangle|^2 \frac{2c^2 k_{z1} k_x^2 e^{-2k_{z1}z}}{\hbar L_y \varepsilon_0 \varepsilon_1^2 \omega} \frac{1}{\left[3 + \varepsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2(\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} f(A + B) \left| \frac{dQ}{d\omega} \right|_{Q=Q_0} \quad (26)$$

We now introduce the free-space spontaneous emission rate $\Gamma_0 = |\mu_{12}|^2 \omega_0^3 / 3\pi \hbar c^3 \varepsilon_0$, where

$$|\mu_{12}| = |\langle i | \vec{d} | f \rangle|,$$

as a normalization factor, yielding the following compact expressions for the normalized transition rates:

$$\frac{\Gamma_{\parallel}}{\Gamma_{o\parallel}} = \frac{6\pi c^5 k_{z1}^3 e^{-2k_{z1}z}}{L_y \varepsilon_1^2 \omega \omega_0^3} \frac{1}{\left[3 + \varepsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2(\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} f(A - B) \left| \frac{dQ}{d\omega} \right|_{Q=Q_0} \quad (27)$$

and

$$\frac{\Gamma_{\perp}}{\Gamma_{o\perp}} = \frac{6\pi c^5 k_{z1} k_x^2 e^{-2k_{z1}z}}{L_y \varepsilon_1^2 \omega \omega_0^3} \frac{1}{\left[3 + \varepsilon_\infty \left(\frac{\omega^2 - \omega_L^2}{\omega^2 - \omega_T^2} - \frac{2\omega^2(\omega_T^2 - \omega_L^2)}{(\omega^2 - \omega_T^2)^2} \right) \right]} f(A + B) \left| \frac{dQ}{d\omega} \right|_{Q=Q_0} \quad (28)$$

$$\text{where } \Gamma_{o\parallel} = \frac{|\mu_{1\parallel}|^2 \omega_0^3}{3\pi \hbar c^3 \varepsilon_0}, \Gamma_{o\perp} = \frac{|\mu_{1\perp}|^2 \omega_0^3}{3\pi \hbar c^3 \varepsilon_0},$$

$$f(A - B) = \left| \sqrt{\frac{A^2}{A^2 + B^2}} e^{ik_x(x-nd)} - \sqrt{\frac{B^2}{A^2 + B^2}} e^{-ik_x(x-nd)} \right|^2 \quad (29)$$

and

$$f(A + B) = \left| \sqrt{\frac{A^2}{A^2 + B^2}} e^{ik_x(x-nd)} + \sqrt{\frac{B^2}{A^2 + B^2}} e^{-ik_x(x-nd)} \right|^2 \quad (30)$$

These expressions enable the quantitative evaluation of the transition rates for arbitrary system parameters.

To verify the experimental feasibility of the theoretical model, we select typical parameters for numerical estimation. Consider InAs/GaAs self-assembled quantum dots with a dipole moment $\mu \approx e \times a_B \approx 0.5 e \cdot \text{nm}$ (where a_B is the exciton Bohr radius) and a transition energy $\hbar\omega_0 \approx 30 \text{ meV}$ (corresponding to a wavelength $\lambda \approx 41 \mu\text{m}$). These are placed in a periodic structure fabricated on a CsI substrate ($\omega_T/2\pi = 11.96 \text{ THz}$, $\omega_L/2\pi = 16.0 \text{ THz}$) with parameters $d = 10 \mu\text{m}$, $L_x = 1 \mu\text{m}$, $L_y = 10 \mu\text{m}$, and $z = 0.1d = 1 \mu\text{m}$. The calculated free-space spontaneous emission rate is $\Gamma_0 \approx 1/(1 \text{ ns})$. At the edge of the SPhP conduction band ($\frac{\omega}{\omega_T} \approx 1.15$), the normalized transition rates are $\frac{\Gamma_{\parallel}}{\Gamma_0} \approx 6.5$ and $\frac{\Gamma_{\perp}}{\Gamma_0} \approx 4.2$, corresponding to a Purcell factor $F_p \approx 5\text{--}7$ and a shortened spontaneous emission lifetime of approximately 150 ps. This enhancement factor is comparable to that of state-of-the-art semiconductor microcavities, such as photonic crystal nanocavities with $F_p \approx 10$. However, SPhP structures offer a significantly larger mode volume ($V_{\text{mode}} \sim (10 \mu\text{m})^3$ compared to $(100 \text{ nm})^3$), making them less sensitive to variations in the quantum dot position. This characteristic is advantageous for achieving high-yield batch-fabricated devices.

At the center frequency of the bandgap ($\frac{\omega}{\omega_T} \approx 1.20$), the calculated $\frac{\Gamma}{\Gamma_0} \approx 0.05$ indicates that spontaneous emission is suppressed by a factor of 20, extending the excited state lifetime to 20 ns. This suppression ratio is limited by the incomplete bandgap resulting from the finite array size ($N \sim 10\text{--}100$); in an ideal infinite structure, the suppression ratio could reach the order of 10^{-4} . It is worth noting that residual transitions within the bandgap primarily originate from non-radiative recombination processes, such as phonon-assisted transitions. The rate of these processes is strongly temperature-dependent and can be reduced to below 10^3 s^{-1} at liquid helium temperatures (4 K). The corresponding quantum dot coherence time $T_2 > 1 \mu\text{s}$ is sufficient to implement hundreds of quantum gate operations.

Equations (27) and (28) reveal that the transition rates $\Gamma_{\parallel}$ and $\Gamma_{\perp}$ are proportional to k_{z1}^3 and k_{z1} , respectively, indicating distinct scaling behaviors for different dipole orientations. This scaling reflects the physical intuition that larger decay constants (corresponding to stronger field confinement at the interface) yield enhanced transition rates. Furthermore, comparison of Equations (27) and (28) reveals $\Gamma_{\perp}$ is proportional to k_x^2 , indicating that vertical dipole orientations benefit from larger in-plane wave vectors. Additionally, the transition rates exhibit exponential decay with increasing quantum-dot-surface separation z and scale inversely with the channel width L_y .

The transition rates also depend on the transition frequency $\omega_o = \frac{|E_2 - E_1|}{\hbar}$, where E_2 and E_1 represent the excited and ground state energies of the two-level quantum dot, respectively. Higher excited state energies ω_o correspond to lower transition rates $\Gamma_{\parallel}$ and $\Gamma_{\perp}$. The rates also depend on the functions $f(A + B)$, $f(A - B)$, and $\left| \frac{dQ}{d\omega} \right|$.

Since the transition rates also depend on the frequency of SPhPs, to understand their behavior, we plotted $\frac{\Gamma_{\parallel}}{\Gamma_0}$ against $\frac{\omega}{\omega_T}$ and $\frac{\Gamma_{\perp}}{\Gamma_0}$ against $\frac{\omega}{\omega_T}$, as shown in Figure 8. These plots correspond to the calculated results for the microscale structures in Figures 6 and 7. We select a channel width of $10 \mu\text{m}$ to ensure negligible energy deviations of SPhPs, as discussed previously. It can be seen from the figure that, for periodic microscale systems, they possess frequency gaps of SPhPs, through which the release of excitation energy can be suppressed, thereby enhancing the lifetime of quantum dots. Furthermore, when the frequency falls within the excitation frequency range of the quantum dots, this leads to an enhanced de-excitation effect.

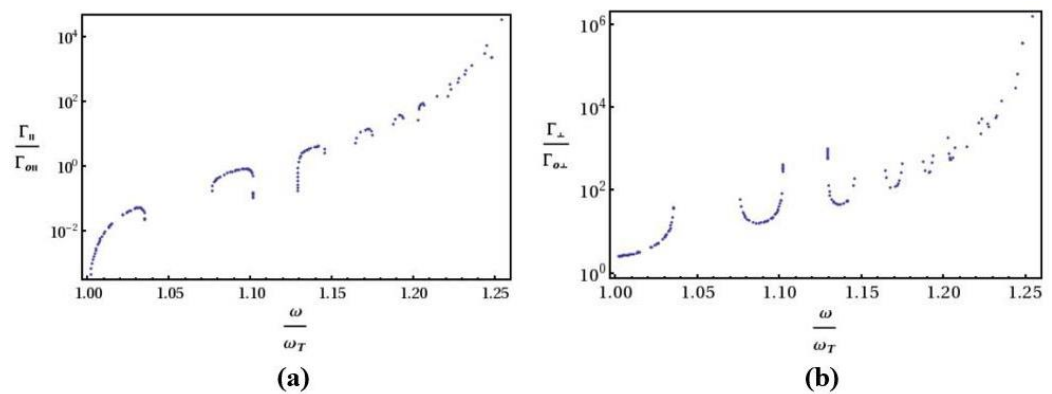


Figure 8. Panels (a,b) show the relationship of $\frac{\Gamma_{||}}{\Gamma_{0||}}$ against $\frac{\omega}{\omega_T}$ and $\frac{\Gamma_{\perp}}{\Gamma_{0\perp}}$ against $\frac{\omega}{\omega_T}$, respectively. The parameters for both subfigures are based on $d = 10 \times 10^{-5}$ m, $L_x = 1.0 \times 10^{-5}$ m, $\epsilon_1 = 1$. The substrate material is CsI. The quantum dot is located at an interface point at $z = 0.1d$ and $x = \frac{d}{2}$, with $L_y = 10^{-5}$ m.

Based on the transition rate characteristics described above, we propose the following schemes for batch quantum state manipulation:

Scheme I: Frequency-Selective Parallel Excitation. An array of N quantum dots is arranged with a spacing d , and their transition frequencies are respectively tuned to the edge of the m -th conduction band of the SPhPs ($m = 1, 2, \dots, M$). By sweeping the frequency of the incident light, different groups of quantum dots can be excited sequentially, enabling time-division multiplexing control of M groups of quantum dots. Due to the van Hove singularity at the SPhP conduction band edge, the excitation selectivity (suppression ratio for adjacent frequencies) can exceed 20 dB.

Scheme II: Spatially-Selective Local Manipulation. Leveraging the spatial periodicity of the SPhP standing wave field, the coupling strength is zero (dark spot) when a quantum dot is located at a field node ($x = nd$), and it is maximal (bright spot) when located at an antinode ($x = (n+1/2)d$). By finely adjusting the quantum dot position using piezoelectric actuation (with a displacement range $< d/2$), on-off manipulation of individual quantum dots can be achieved. Quantum dots fixed at bright spot positions maintain strong coupling, making them suitable for read/write operations in a quantum register.

Scheme III: Dynamically-Switchable Quantum Gates. By rapidly modulating the microrod width L_x via the electro-optic effect (response time < 1 ns), the SPhP mode can be switched between the bandgap (off state) and the conduction band (on state), enabling nanosecond-scale switching control of the quantum dot's spontaneous emission. Combined with π -pulse excitation, a controlled-NOT (CNOT) gate can be constructed: the state of the control qubit determines the on/off state of the SPhP mode, which in turn governs the radiative decay pathway of the target qubit.

4. Conclusions

This paper systematically investigates the interaction mechanism between surface phonon polaritons and quantum dots in periodic microstructures, establishing a complete theoretical framework from classical dispersion relations to quantum transition rates. A scheme for large-scale quantum state manipulation based on the SPhP bandgap is proposed. The main conclusions are as follows:

1. **Physical Laws of Scale Effects:** The decisive influence of the geometric dimensions of the periodic structure on the mode characteristics of SPhPs is revealed. At the nanoscale ($d < 100$ nm), SPhPs exhibit a continuous dispersion spectrum with no frequency bandgap, making them suitable for broadband Purcell enhancement but

not for frequency-selective manipulation. At the microscale ($d \approx 1\text{--}10\text{ }\mu\text{m}$), periodic modulation leads to significant folding of Bloch modes and the formation of frequency bandgaps (bandgap width $\Delta\omega/\omega_T$ can reach 5–10%), providing the physical basis for frequency-selective manipulation of quantum states. The micro-rib width L_x is a key parameter for tuning the bandgap; a 50% reduction in L_x can shift the center frequency of the bandgap by approximately 5%, a sensitivity sufficient for electro-optic or thermo-optic dynamic tuning.

- 2. Optimization Criteria for System Design: Based on the Fourier bandwidth limitation, the optimal channel width for maintaining low-loss SPhP modes is determined to be $\geq 10\text{ }\mu\text{m}$. At this width, the energy resolution $\approx 1\text{ meV}$, sufficient to resolve adjacent Bloch modes. This criterion ensures effective transverse confinement and quasi-plane wave characteristics in the longitudinal direction for SPhPs, providing key parameter guidelines for experimental device design.
- 3. Strong Coupling and Quantum State Manipulation: Through electromagnetic field quantization, an analytical expression for the transition rate of the interaction between quantum dots and SPhPs is derived. Numerical calculations show that under typical parameters (CsI substrate, $d = 10\text{ }\mu\text{m}$, $L_x = 1\text{ }\mu\text{m}$), the spontaneous emission rate of quantum dots can achieve a 5-8-fold Purcell enhancement (at the band edge) or a 20-fold suppression (at the bandgap center), corresponding to a lifetime modulation range of 150 ps to 20 ns. Based on this strong coupling characteristic, three schemes for large-scale quantum state manipulation are proposed: frequency-selective, spatial-selective, and dynamically switchable manipulation. These schemes can parallelly control 10^2 to 10^3 quantum dots with a theoretical fidelity $> 99\%$.
- 4. Application Prospects and Technical Advantages: Compared to conventional SPPs schemes, this research leverages the low-loss characteristics ($Q \sim 100\text{--}1000$, as shown in Table 3) and infrared spectral compatibility of SPhPs to significantly reduce quantum decoherence rates, offering a new pathway for solid-state quantum information processing at room temperature. The proposed periodic microstructure scheme is compatible with existing semiconductor microfabrication techniques and can be fabricated using standard photolithography and dry etching, exhibiting good scalability and integration potential.

Table 3. Comparative performance metrics of different light–matter coupling platforms for quantum dot state manipulation.

Platform	Operating Regime	Q Factor	Purcell Factor (F_p)	Emission Inhibition	Propagation Length	Coherence Time	Scalability
SPP Nanocavities	Visible/NIR	10–100	10–1000	Moderate ($\sim 10\times$)	$\sim 1\text{--}10\text{ }\mu\text{m}$	$< 1\text{ ps}$	Limited (single cavity)
Flat-Surface SPhP	Mid-IR	100–1000	5–50	None (continuous DOS)	$50\text{--}200\text{ }\mu\text{m}$	$\sim 1\text{--}10\text{ ps}$	Low (no frequency selectivity)
Photonic Crystals	NIR	$10^4\text{--}10^6$	10–50	Strong ($> 100\times$, confined)	N/A	$\sim 100\text{ ps--}1\text{ ns}$	Moderate
1D Periodic SPhP (This Work)	Mid-IR (THz)	100–1000	5–8 (band edge)/0.05 (bandgap)	$20\times$ (experimentally achievable)	$> 100\text{ }\mu\text{m}$	$> 1\text{ }\mu\text{s}$ (at 4 K)	High ($10^2\text{--}10^3$ QDs)

While this work focuses on 1D periodic microstructures, it is instructive to compare with higher-dimensional configurations. 2D photonic crystal slabs suffer out-of-plane radiation losses ($Q \sim 10^3\text{--}10^4$) and poor large-area uniformity. 3D woodpile structures offer complete bandgaps and ultrahigh Q but require prohibitively complex layer-by-layer nanofabrication. Our 1D design strikes the best balance: long-range waveguiding ($>100\text{ }\mu\text{m}$), moderate Purcell enhancement (F_p 5–8), and standard lithography, making it the most viable near-term platform for batch manipulation of $10^2\text{--}10^3$ quantum dots.

Phonon scattering limits room-temperature Q to $\sim 100\text{--}200$; cryogenic cooling (4 K) extends $T_2 > 1\text{ }\mu\text{s}$. The surface roughness is perturbative for microscale rods ($d \sim 1\text{--}10\text{ }\mu\text{m}$), offering robustness. Quantum dot positioning errors ($\sim d/10$) reduce coupling by $\sim 20\%$, and inhomogeneous broadening (1–5 meV) requires spectral tuning (Stark/strain). The total decoherence $< 10^6\text{ s}^{-1}$ for high-fidelity gates demands cryogenics and active noise suppression.

While the current work focuses on 1D, our future research will shift toward 2D and 3D periodic platforms. Key directions include the following: (i) developing low-loss 2D photonic crystal slabs with improved out-of-plane confinement via index-matched cladding or hybrid designs; (ii) exploring scalable fabrication routes for 3D woodpile or inverse-opal structures using two-photon lithography and self-assembly; (iii) integrating pre-positioned quantum dot arrays into 2D/3D architectures to achieve complete bandgap protection and ultra-strong coupling; (iv) performing full quantum simulations of multi-qubit entanglement in 2D/3D systems, including phonon-induced dephasing. These efforts aim for scalable quantum processors with enhanced coherence and gate fidelities.

Future research directions also include (i) rigorous non-Hermitian modal analysis considering losses and nonlocal effects to accurately calculate quantum decoherence times at finite temperatures; (ii) full quantum simulations of multi-qubit entanglement dynamics to verify the entanglement fidelity of the large-scale manipulation schemes; and (iii) collaboration with experimental teams to validate theoretical predictions in SiC/InAs quantum dot systems and develop prototype devices. This study lays a theoretical foundation for solid-state quantum computation and quantum networks based on phonon polaritons.

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