

Efficient coupling of light to an atomic tweezer array in a cavity

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(Received 25 January 2024; accepted 4 December 2024; published 30 December 2024)

We consider the coupling of light, via an optical cavity, to two-dimensional atomic arrays whose lattice spacing exceeds the wavelength of the light. Such “superwavelength” spacing is typical of optical tweezer arrays. While subwavelength arrays exhibit strong atom-photon coupling, characterized by high optical reflectivity in free space, the coupling efficiency of superwavelength arrays is reduced due to collective scattering losses to high diffraction orders. We show how a moderate-finesse cavity can overcome these losses. As the scattering losses peak at certain discrete values of the lattice spacing, the spacing can be optimized to achieve efficient atom-photon coupling in the cavity. Our cavity-QED theory properly accounts for collective dipolar interactions mediated by the lossy, non-cavity-confined photon modes and for finite-size effects of both the array and the light field. These findings pave the way to harnessing the versatility of tweezer arrays for efficient atom-photon interfaces in applications of quantum computing, networking, and nonlinear optics.

DOI: 10.1103/PhysRevResearch.6.L042070

Introduction. Spatially ordered arrays of atoms trapped by optical tweezers have emerged as a notable platform for quantum science and technology [1–14]. Here, we discuss the promising potential of establishing an efficient quantum interface between such atomic arrays and light. The motivation is twofold. First, viewing tweezer arrays as stationary nodes of quantum information systems, their coupling to light is essential for quantum networking or for distributed quantum computing [15–18]. Second, the nonlinearity and versatile tools developed for atomic arrays are useful for generating and studying quantum many-body states of light [19–23].

Atom-photon interfaces can be characterized by the coupling efficiency $r_0 = \Gamma/(\Gamma + \gamma_{\text{loss}})$, with Γ the coupling rate of the atomic system to a desired “target” mode of the photons and γ_{loss} the loss rate from scattering to other modes [24]. Importantly, in atomic arrays, the latter is dominated by collective effects established by induced dipole-dipole interactions between the atoms. For subwavelength arrays, where the array lattice spacing a is smaller than the wavelength of light λ , typical of optical lattices [25–27], such collective effects significantly reduce the scattering losses γ_{loss} , leading to efficient coupling characterized by the array’s strong reflectivity [28–48]. Conversely, in a superwavelength lattice $a > \lambda$, typical of tweezer arrays [4,5], the atoms collectively couple not only to the normal-incident target mode but also to higher diffraction orders [24,30]. These diffraction orders translate to scattering losses γ_{loss} , resulting in drastically reduced reflectivity and efficiencies of quantum operations [24].

Here, we consider the prospect of overcoming these losses by using an optical cavity. The idea, common to various atomic systems, is that multiple reflections of the target mode between the cavity mirrors enhance its coupling Γ to atoms, while leaving unaffected the losses γ_{loss} from off-cavity axis scattering. However, achieving efficient coupling, $\Gamma/\gamma_{\text{loss}} \gg 1$, requires also to properly analyze the losses γ_{loss} . Since in arrays the losses originate in collective scattering, the analysis must go beyond typical cavity QED theory which neglects collective effects due to modes outside of the cavity [49,50].

To this end, we study a two-dimensional (2D) superwavelength array placed inside a cavity, while properly accounting for collective dipole-dipole interactions mediated by non-cavity-confined modes. First, this allows to find the cavity enhancement of the desired coupling $\Gamma \propto \mathcal{F}$ by the cavity

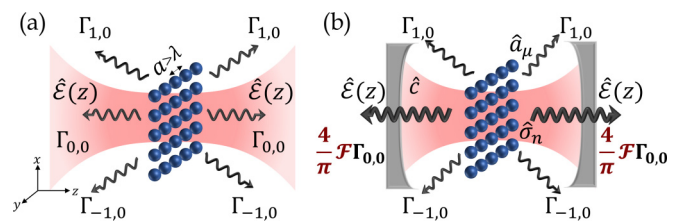


FIG. 1. (a) Coupling light to a 2D array in free space. The array with lattice spacing a is located on the xy plane and illuminated from both sides by a paraxial target mode $\hat{\mathcal{E}}(z)$, corresponding to the $(0,0)$ diffraction order. In the superwavelength case $a > \lambda$, scattering also exists to higher diffraction orders (m_x, m_y) , thus contributing to losses [Eq. (6)]. Here, the orders $(-1, 0)$ and $(1, 0)$ are marked, and $\Gamma_{0,0} \equiv \Gamma_0$ [Eq. (5)]. (b) A 2D array inside an optical cavity. The coupling to the $(0,0)$ order is enhanced by the cavity finesse $\mathcal{F}$, while the coupling to all other diffraction orders is unaffected.

fineness $\mathcal{F}$, thus generalizing this principle to collective array systems. More importantly, our treatment of collective effects proves essential for minimizing the losses γ_{loss} . We find that the losses are dominated by the collective scattering to higher diffraction orders outside the cavity, even when losses due to finite-size effects and imperfections are present. Crucially, these dominant collective losses depend strongly on the array lattice spacing a , with large losses appearing at specific resonant values of a , accompanied by nontrivial dependence on the cavity-mode waist. Accordingly, our analytical results, supported by numerical scattering calculations, reveal the possibility of designing near-unity coupling efficiencies r_0 by a proper choice of the lattice constant, even for moderate fineness. This opens the way for achieving high fidelities of various quantum tasks performed by the tweezer-array quantum interface.

General formalism. We consider a 2D tweezer array consisting of N identical two-level atoms, forming a square lattice on the xy plane with a superwavelength lattice spacing $a > \lambda$, with λ being the central wavelength of the incident light. In the following, we will first consider an array in free space [Fig. 1(a)] and then proceed to the case where it is placed inside a cavity [Fig. 1(b)]. In both cases, we will study how the array is coupled to a paraxial “target mode” of interest $\hat{\mathcal{E}}(z) = \frac{1}{\sqrt{2}}[\hat{\mathcal{E}}_+(z) + \hat{\mathcal{E}}_-(-z)]$, illuminating the array from both sides ($\pm$ for right- and left-propagating beams, respectively) [51]. To this end, we follow the approach of Ref. [24] wherein many-atom, multimode problems are mapped to a one-dimensional (1D) scattering problem, of a target mode $\hat{\mathcal{E}}$ with a transverse mode profile $u(\mathbf{r}_\perp)$ [$\mathbf{r}_\perp = (x, y)$] coupled to a spatially matched collective-dipole $\hat{P}$,

$$\hat{P} = \frac{a}{\sqrt{\eta}} \sum_{n=1}^N u^*(\mathbf{r}_n^\perp) \hat{\sigma}_n. \quad (1)$$

Here, $\hat{\sigma}_n$ is the two-level lowering operator of atom n , and

$$\eta = \int_{L_a^2} |u(\mathbf{r}_\perp)|^2 d\mathbf{r}_\perp \xrightarrow{\text{Gaussian}} \eta = \text{erf}^2\left(\frac{L_a}{\sqrt{2}w}\right), \quad (2)$$

is the fraction of the spatial mode that overlaps with the atomic array of size $L_a^2 = a^2 N$. For a typical Gaussian target mode, $u(\mathbf{r}_\perp) = \sqrt{\frac{2}{\pi w^2}} e^{-(\mathbf{r}_\perp)^2/w^2}$ with waist $w \gg \lambda$, excellent overlap $\eta \rightarrow 1$ exists for a sufficiently large array $L_a \gg w$.

Within the general 1D description, the coupling between the target mode and the collective dipole in the low-excitation, linear regime (where $\hat{P}$ is bosonic, $[\hat{P}, \hat{P}^\dagger] = 1$) is captured by the Heisenberg equations [24]

$$\begin{aligned} \frac{d\hat{P}}{dt} &= -\left[\frac{\Gamma}{2} + \frac{\gamma_{\text{loss}}}{2} + i(\Delta - \delta)\right]\hat{P} + i\sqrt{\Gamma}\hat{\mathcal{E}}_0(0) + \hat{F}, \\ \hat{\mathcal{E}}(z) &= \hat{\mathcal{E}}_0(z) + i\sqrt{\Gamma}\hat{P}. \end{aligned} \quad (3)$$

Here δ is the detuning between the central frequency of the target mode and the bare atomic transition, Γ is the coupling to the target mode (whose input field is $\hat{\mathcal{E}}_0$), γ_{loss} is the loss rate to undesired modes, Δ is a collective resonance shift, and $\hat{F}$ is the quantum Langevin noise corresponding to γ_{loss} . The quality of the array-light interface in the 1D model is fully characterized by the cooperativity C , or equivalently, by the

coupling efficiency r_0 , defined as

$$C = \frac{\Gamma}{\gamma_{\text{loss}}}, \quad r_0 = \frac{C}{1+C}. \quad (4)$$

In a two-sided system, where light can either be transmitted or reflected, r_0 is also equal to the on-resonance reflectivity. These parameters determine the fidelity of various quantum tasks, e.g., the quantum memory efficiency, or the maximal level of quantum correlations achievable when adding nonlinearity [24]. We thus require $\Gamma \gg \gamma_{\text{loss}}$ for an efficient light-matter coupling.

Array in free space. To understand the sources of scattering losses, it is instructive to consider first an array in free space. The mapping of the many-atom, multimode dynamical equations to those of Eqs. (3) is performed by assuming the following: (i) the array is large enough, $L_a \gg \lambda$, to take collective effects due to dipole-dipole interactions as in an infinite array; (ii) the target mode is paraxial, $w \gg \lambda$, neglecting spatial dispersion effects in the array response (corrections are discussed below). With these assumptions, we obtain Eqs. (3) with the parameters [24,30]

$$\begin{aligned} \Gamma^{(\text{free})} &= \eta\Gamma_0 = \eta \frac{3}{4\pi} \frac{\lambda^2}{a^2} \gamma, \\ \gamma_{\text{loss}}^{(\text{free})} &= \gamma_s + (1-\eta)\Gamma_0 + \gamma_{\text{diff}}, \end{aligned} \quad (5)$$

and $C^{(\text{free})} = \Gamma^{(\text{free})}/\gamma_{\text{loss}}^{(\text{free})}$. Here Γ_0 represents the collective emission rate of a uniformly excited infinite array and γ is the free-space spontaneous emission rate of a single atom.

The three terms comprising $\gamma_{\text{loss}}^{(\text{free})}$ originate from distinct loss mechanisms. The first two terms are associated with imperfections, with γ_s being the scattering loss due to noncollective processes such as weak disorder [24], and $(1-\eta)\Gamma_0$ representing the loss due to imperfect overlap of the target mode with the array (with $\gamma_s \ll \Gamma_0$ assumed negligible in the following unless stated otherwise). The third term γ_{diff} exists only in a superwavelength lattice ($a > \lambda$), and it describes the emission to higher diffraction orders

$$\gamma_{\text{diff}} = \Gamma_0 \sum_{m_x, m_y} \frac{1 - \frac{\lambda^2}{a^2} |(m_x, m_y) \cdot \mathbf{e}_d|^2}{\sqrt{1 - \frac{\lambda^2}{a^2} (m_x^2 + m_y^2)}} \equiv \sum_{m_x, m_y} \Gamma_{m_x, m_y}. \quad (6)$$

Here, the sum $\sum'$ is preformed over all propagating diffraction orders $(m_x, m_y) \neq (0, 0)$ satisfying the condition $|(m_x, m_y)| < a/\lambda$.

In an imperfect subwavelength array, enhancing the coupling Γ to the target mode using a cavity is of course beneficial for increasing $C = \Gamma/\gamma_{\text{loss}}$. However, for a superwavelength array, it becomes practically *essential* even in a perfect infinite system due to the scattering $\gamma_{\text{diff}} > 0$ to multiple diffraction orders which always exists and often dominates the loss, as seen in Fig. 2 and discussed further below.

Array inside an optical cavity. Motivated by the above conclusions, we proceed to consider an array inside a cavity. We assume that the target mode spatially matches the cavity mode so that its coupling to the array is enhanced, whereas the additional diffraction orders are all directed at angels which are not supported by the cavity mirrors and are thus scattered away, see Fig. 1(b). Since the scattering to the high diffraction orders is collective, the model must take into

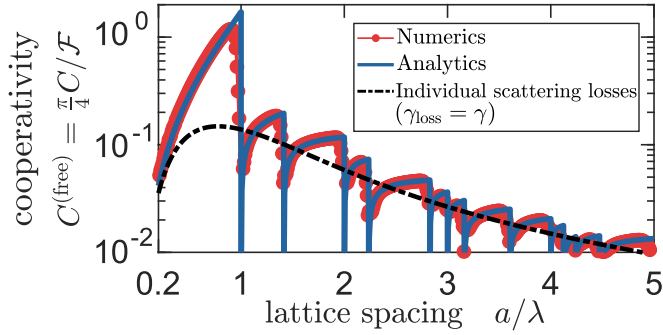


FIG. 2. Cooperativity of an atomic array in free space $C^{(\text{free})}$ and in an optical cavity C , noting their equivalence up to a factor proportional to the cavity finesse $C^{(\text{free})} = \frac{\pi}{4} C/\mathcal{F}$ [Eq. (8)]. Analytical results [solid blue, Eqs. (5) and (6)] and numerical scattering results (red circles) are obtained for an array of $N = 20 \times 20$ atoms and a beam waist $w = 15\lambda$. The cooperativity changes rapidly around the points $a/\lambda = \{1, \sqrt{2}, 2, \sqrt{5}, \dots\}$, where the new diffraction orders appear. These diffraction resonances are absent when collective effects in the scattering losses are neglected (dashed black, $\gamma_{\text{loss}} = \gamma$).

account dipole-dipole interactions mediated not only by the cavity modes but also by the nonconfined three-dimensional (3D) modes. To this end, we develop a modified cavity QED formalism that accounts for such collective effects, adopting the approach recently proposed in Ref. [53] to the case of a superwavelength array.

We begin with the Hamiltonian of the full many-atom multimode system depicted in Fig. 1(b) (see Ref. [52]). The atoms $\hat{\sigma}_n$ are coupled to the cavity mode $\hat{c}$ via the Hamiltonian $\hat{H}_{AC} = \hbar \sum_n [g_n \hat{\sigma}_n^\dagger + \text{H.c.}]$, with $g_n \propto u(\mathbf{r}_n^\perp)$. In turn, the cavity mode is coupled with a damping rate κ to the 1D continuum representing the propagating target mode $\hat{\mathcal{E}}$. In addition, the atoms are coupled to a photon reservoir comprised of the modes $\hat{a}_\mu$ which are not confined by the cavity. As detailed in Ref. [52], we first write Heisenberg-Langevin equations for the atomic and cavity variables by eliminating the nonconfined modes $\hat{a}_\mu$ within a Markov approximation. For the equations of the atomic variables $\hat{\sigma}_n$, the nonconfined modes contribute a collective interaction term between dipoles of different atoms, with the dipole-dipole kernel given by a partial photonic Green's function G_{nc} , which contains only the nonconfined modes. In general, finding the nonconfined modes and G_{nc} is a system-specific, nontrivial task. However, a generic approximation can be made noting that (i) the dispersive shift (given by $\text{Re}[G_{\text{nc}}]$) is dominated by near fields and is hence similar to free space, assuming the array is at least a wavelength away from the cavity mirrors and (ii) the emission to zeroth diffraction order is into the cavity, and hence must not be included in the losses due to the nonconfined modes (given by the $\text{Im}[G_{\text{nc}}]$). These considerations are accounted for by the following approximation [52–54]:

$$\text{Re}[G_{\text{nc}}] \simeq \text{Re}[G], \quad \text{Im}[G_{\text{nc}}] \simeq \text{Im}[G] - \text{Im}[G_c]. \quad (7)$$

Here G is the free-space Green's function of the photons, and G_c is the Green's function of the 1D continuum corresponding to the transverse cavity mode, which does not contain the

modes corresponding to the additional diffraction orders that diverge aside.

With this approximation and a subsequent adiabatic elimination of the fast decaying cavity mode $\hat{c}$, we show in Ref. [52] that we arrive at the equations in the form of Eq. (3), with the parameters

$$\Gamma = \frac{4}{\pi} \mathcal{F} \Gamma^{(\text{free})}, \quad \gamma_{\text{loss}} = \gamma_{\text{loss}}^{(\text{free})}, \quad C = \frac{\Gamma}{\gamma_{\text{loss}}} = \frac{4}{\pi} \mathcal{F} C^{(\text{free})}. \quad (8)$$

Namely, compared to the free-space case, the coupling to the target mode obtains an enhancement factor proportional to the cavity finesse $\mathcal{F}$, while the scattering losses remain the same since the higher diffraction orders are spanned by the nonconfined modes described by G_{nc} in Eq. (7).

The enhancement of the cooperativity C by the finesse generalizes known results of noncollective behavior in a cavity, e.g., a single atom or a dilute ensemble [49,50,55] to the collective case of an array. Indeed, considering a reference case of an atomic ensemble with density ρ and length L_z placed in a cavity, and neglecting collective dipolar interactions via nonconfined modes as in typical cavity QED theory [49,50] (justified for dilute ensembles), one obtains the individual-decay loss rate $\gamma_{\text{loss}}^{(\text{ref})} = \gamma$ and a coupling $\Gamma^{(\text{ref})}$ in the form of Eq. (8) but with the average transverse distance $a_{\text{ref}} = 1/\sqrt{\rho L_z}$ replacing the lattice spacing a [52]. Namely, while a similar cavity enhancement is obtained for Γ , the distinctive feature of our theory is its collective treatment of losses, which turns out to be crucial, as discussed below.

Diffraction resonances. With Eq. (8) at hand [alongside Eqs. (5) and (6)], we now examine the cooperativity C as a function of the lattice spacing, for a given finite atom number N and target-mode waist w . The result is plotted in Fig. 2, wherein, following Eq. (8), C is scaled by $(4/\pi)\mathcal{F}$ highlighting its connection to the free-space result $C^{(\text{free})}$. We validate our analytical results by using numerical calculations of classical scattering in the simpler case of an array in free space illuminated by a Gaussian beam with the same waist. There, $C^{(\text{free})}$ is obtained from the efficiency $\eta_0^{(\text{free})} = C^{(\text{free})}/(1 + C^{(\text{free})})$, which we recall to be given by the on-resonance reflectivity estimated from the scattered field as described in Refs. [24,30], finding excellent agreement with the analytical results [see Fig. 2; sharper features in the analytical curve are due to the infinite-array approximation in Eq. (6) for γ_{diff}].

We observe that the cooperativity C in Fig. 2 does not change smoothly with the lattice spacing a but rather drops rapidly to zero at specific values. These values occur at $a/\lambda = \{1, \sqrt{2}, 2, \sqrt{5}, \dots\}$ where the diffraction losses in Eq. (6) exhibit divergences, corresponding to the emergence of new dissipation channels: diffraction orders that change from being evanescent (nondissipative) to propagating (radiative). Enhanced cooperativity is obtained for lattice spacings a just before new diffraction orders appear, corresponding to minima of γ_{diff} in Eq. (6). This nontrivial behavior is in stark contrast to the reference case, the prediction of a typical cavity QED theory that neglects collective effects in the emission to nonconfined modes [49] (dashed curve in Fig. 2). Our collective theory is, therefore, essential for designing efficient

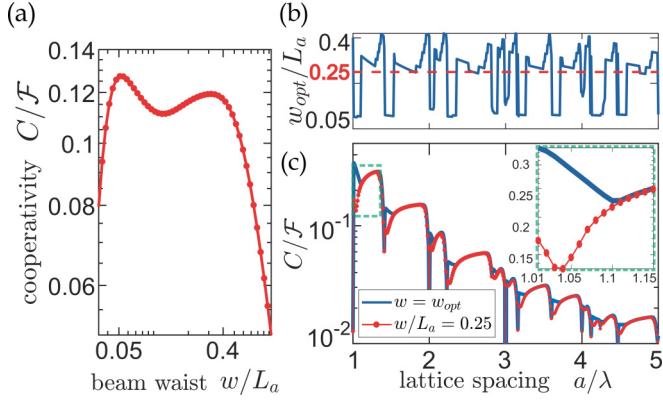


FIG. 3. Finite-size effects for an array of $N = 20 \times 20$ atoms in a cavity. (a) Cooperativity versus the beam waist w close to the diffraction resonance $a/\lambda = \sqrt{2} + 0.1$. (b) The optimal waist versus the lattice spacing. The optimal waist is close to the average value $w/L_a = 0.25$, besides the jumps around the diffraction resonances. (c) The cooperativities for the optimal waists from (b), and for the average value $w/L_a = 0.25$ become distinct closer to the diffraction resonances (inset). The results are obtained via numerical scattering calculations using $C^{(free)} = \frac{\pi}{4} C/F$.

interfaces, particularly by choosing the optimal values of a just before the resonances.

Finite-size effects. We now show how the diffraction resonances also lead to nontrivial behavior of the optimal waist. So far, the finite size of the array and beam is taken into account via the error function η . While this is an excellent approximation for a paraxial light-beam, it clearly favors small beams with respect to the array size. However, when optimizing the beam size, one has to take into account a competing *dispersion effect*. Namely, as the beam waist narrows, nonzero transverse momentum $\mathbf{k}_\perp$ components, which correspond to different collective dipoles of the array with different resonances $\Gamma_{\mathbf{k}_\perp}/2 + i\Delta_{\mathbf{k}_\perp}$, become more significant, and the frequency of the beam cannot match all $\Delta_{\mathbf{k}_\perp}$ simultaneously, resulting in reduced cooperativity [30,31].

In superwavelength arrays, the physics is richer: an additional *diffraction effect* occurs around the diffraction resonances. As discussed earlier, when the lattice spacing crosses a point of a new diffraction order, determined by Eq. (6), the cooperativity drops rapidly to zero. However, for nonzero transverse momentum components $\mathbf{k}_\perp$, the diffraction orders appear at different points since, for them the wave vector $\frac{1}{a}(m_x, m_y)$ in Eq. (6) changes to $\frac{1}{a}(m_x, m_y) + \mathbf{k}_\perp/(2\pi)$. Therefore, a smaller waist with a larger momentum spread could reduce the rapid change in the cooperativity. As a result, the dependence of C on the waist may exhibit two maxima. The first maximum results from the balance between the dispersion effect, which favors a larger beam waist, and the diffraction effect, favoring a small beam waist. The second maximum results from the balance between the dispersion effect and the spatial overlap mismatch between the beam and the array. Figure 3(a) shows the cooperativity close to the appearance of the second diffraction order $a/\lambda = \sqrt{2} + 0.1$. Two maxima are seen providing two optimal regimes to work in.

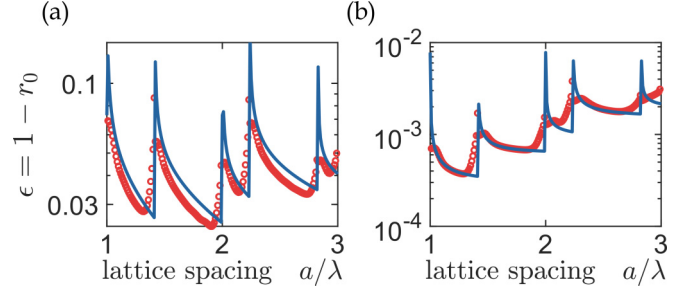


FIG. 4. The coupling inefficiency $\epsilon = 1 - r_0$ for an array of $N = 10 \times 10$ atoms in a cavity, including an individual-atom loss rate $\gamma_s = 0.1\Gamma_0$ accounting for weak disorder or additional decay channels (numerical calculations in red dots; analytical results in solid blue). (a) For cavity finesse $\mathcal{F} = 1000$ and beam waist $w/\lambda = 25$. (b) For $\mathcal{F} = 10000$ and $w/\lambda = 5$. At $a/\lambda \lesssim \sqrt{2}$, the inefficiency reduces to $\epsilon \sim 3.5 \times 10^{-4}$.

Figure 3(b) shows the optimal waist as a function of the lattice spacing. The optimal waist, which is generally close to its average value $w/L = 0.25$, exhibits jumps around the diffraction resonances, corresponding to the two maxima as demonstrated in Fig. 3(a). Figure 3(c) displays the cooperativity for these numerically found optimal waists. We find that around the diffraction resonances, the exact optimal waist gives better results than its average value, while far from these points, they coincide.

Realistic example. The presented approach reveals excellent efficiencies using realistic parameters. Figure 4(a) shows the coupling inefficiency $\epsilon = 1 - r_0 = \frac{1}{1+C}$ for an array of $N = 10 \times 10$ atoms placed inside a cavity with finesse $\mathcal{F} = 1000$ and beam waist $w = 25\lambda$ (for example, $w \approx 20 \mu\text{m}$ for ^{87}Rb atoms with $\lambda = 0.78 \mu\text{m}$). Imperfections such as weak disorder or possible decay to additional stable atomic levels (e.g., in quantum memory schemes) are modeled by the individual-atom loss parameter γ_s in Eq. (5) [24,53,54], taken as $\gamma_s = 0.1\Gamma_0$ in Fig. 4(a) (see Ref. [52] for more examples). It is seen that even for lattice spacing of about $a \approx 1.9\lambda$, the inefficiency is smaller than 2%. Interestingly, since the inefficiency does not decrease monotonically, the inefficiency at $a \approx 1.9\lambda$ is smaller than at the preceding minimum around $a \approx \sqrt{2}\lambda$, demonstrating that sometimes increasing the lattice spacing can reduce the inefficiency. Figure 4(b) shows that a cavity with finesse $\mathcal{F} = 10000$ and a beam waist $w/\lambda = 5$ reduces the inefficiency to $\epsilon \sim 3.5 \times 10^{-4}$.

Discussion. We analyzed the coupling of atomic tweezer arrays to light via an optical cavity. Accounting for collective scattering effects, we revealed an essential principle for achieving efficient coupling, by tuning the lattice spacing near diffraction resonances. These results can be generalized to array structures beyond that of a square lattice, such as a triangular lattice for which we find similar results. Our findings open the door for using tweezer arrays as efficient light-matter interfaces in various quantum tasks and applications: from interconnecting multiple arrays through photons, to quantum memories and quantum-state readout, or quantum nonlinear optics of Rydberg atoms.

Acknowledgments. We acknowledge financial support from the Israel Science Foundation (ISF); the Directorate for Defense Research and Development (DDR&D); the US-Israel Binational Science Foundation (BSF) and US National Science Foundation (NSF); the Minerva Foundation with funding from the Federal German Ministry for Education and

Research; the Center for New Scientists at the Weizmann Institute of Science; the Council for Higher Education (Israel), QUANTERA (PACE-IN); the Helmsley Charitable Trust; and the Estate of Louise Yagour. This research is made possible, in part, by the historic generosity of the Harold Perlman Family.

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