

Oscillation-free unconventional photon blockade from quadrature driving of a Kerr dimer

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We demonstrate unconventional photon blockade in a symmetric Kerr dimer driven with equal-amplitude fields at a 90° phase difference. The minimum intercavity coupling is $J_{\min} = \gamma/4$ at a Kerr nonlinearity $U \ll \gamma$ achievable in standard photonic molecules. The quadrature-driven site emits strongly antibunched light with a smooth, oscillation-free second-order correlator directly resolvable with standard detectors. The scheme operates under continuous-wave and pulsed excitation, and fabrication disorder can be fully compensated by retuning the drive phase and amplitude ratio, removing the need for postfabrication cavity trimming.

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I. INTRODUCTION

Generating single photons on demand is central to quantum communication and quantum information processing [1]. Conventional photon blockade suppresses multiphoton occupation of a resonator by requiring the Kerr nonlinearity U to exceed the cavity decay rate γ [2,3], a phenomenon first observed in cavity QED [4] and subsequently demonstrated in circuit QED [5], and on-chip photonic systems [6]. In practice, the large nonlinearity required restricts conventional photon blockade to discrete emitters—chiefly self-assembled semiconductor quantum dots [7,8] and atoms [4]—where the quantized energy spectrum provides the necessary anharmonicity. Self-assembled quantum dots require individual postfabrication tuning for indistinguishable multisource emission [8], limiting scalability. Extending antibunched light generation to continuous-medium platforms such as quantum-well or two-dimensional (2D) material polaritons [9–11]—where many identical sources can be etched in a single lithographic step—would be transformative, but material nonlinearities in these systems are orders of magnitude below cavity linewidths, placing conventional blockade out of reach.

Unconventional photon blockade (UPB) [12–14] sidesteps this requirement by exploiting destructive quantum interference between two-photon pathways in a coupled cavity dimer, achieving $g^{(2)}(0) \approx 0$ with $U \ll \gamma$. A key limitation is that the interference requires strong intercavity coupling $J \gg \gamma$, inducing rapid oscillations in $g^{(2)}(\tau)$ at frequency $2J$, with an antibunching window $\delta\tau \sim 1/J \ll 1/\gamma$ [15]. At telecom wavelengths this window is only ≈ 100 ps [16], at the edge of

detector resolution. Compounding this, the condition $J \gg \gamma$ is itself challenging in photonic systems: a large coupling with a small loss rate requires extremely high- Q resonators with tight mode overlap, difficult to reconcile with the fabrication tolerances needed for fine-tuned UPB parameters.

Two recent proposals address these constraints. Flayac and Savona [17] proposed replacing the intercavity hopping with a one-way dissipative coupling to suppress the rapid oscillations in $g^{(2)}(\tau)$. While nonreciprocal photonic elements have been demonstrated—for instance, asymmetric backscattering between counterpropagating modes near an exceptional point [18]—the cascaded, unidirectional dissipative coupling between two separate cavities that this scheme requires remains a distinct and demanding requirement. Wang *et al.* [19] engineered cavity networks with single-photon dark states at large γ/J , achieving long-lived photon blockade with $g^{(2)}(\tau) \propto (\gamma\tau)^4$ over a window of $\sim 8/\gamma$.

Bilateral driving of a Kerr dimer with two coherent fields of tunable relative phase has itself been studied as a route to UPB [20–22], with the phase and amplitude ratio scanned to optimize antibunching; a related input-output approach instead mixes the two cavity outputs on a beam splitter to isolate the antibunched mode [15,23], while an emitter-based scheme unifies conventional and unconventional blockade in a two-photon Jaynes-Cummings model [24]. Here we retain the minimal two-cavity geometry but single out the equal-amplitude quadrature drive $F_2 = iF_1$, which produces antibunching directly on a single physical port (site 2) with no output mixing and no nonreciprocal coupling [Fig. 1(a)]. For this drive the coupling threshold is $J_{\min} = \gamma/4$ —a factor of $2\sqrt{2}$ below the single-site value $\gamma/\sqrt{2}$ [13]—and the entire band $\gamma/4 < J < \gamma/\sqrt{2}$, which admits no single-site UPB solution for any U , becomes accessible, including the overdamped regime $J < \gamma/2$ where $g^{(2)}(\tau)$ rises smoothly without the $2J$ oscillations. We give the parameter locus in closed form, showing that on-demand pulsed operation is possible without tailoring the pulse, and that fabrication disorder can be fully compensated by retuning the drive.

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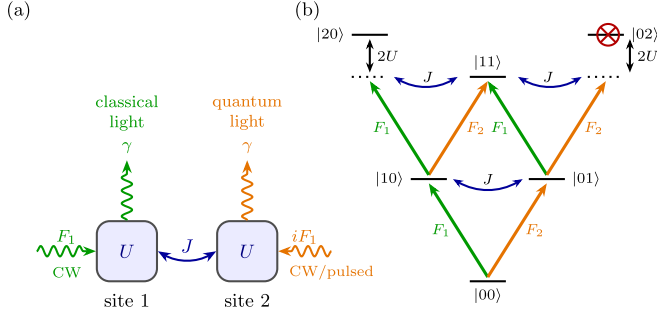


FIG. 1. (a) Kerr dimer setup: site 1 driven by a cw field F_1 and site 2 driven by a pulsed quadrature field iF_1 ; both sites decay at rate γ and are coupled by hopping J . (b) Two-photon energy levels and transition amplitudes. The two-photon occupation of site 2, $|0, 2\rangle$, receives contributions from direct two-photon excitation via F_2 and from indirect transfer via the $|1, 1\rangle$ intermediate state and hopping J . These two channels interfere destructively at the UPB condition, giving zero two-photon occupation of site 2 (crossed circle).

II. MODEL

Two coupled single-mode Kerr cavities are driven coherently under the Hamiltonian (rotating frame, $\hbar = 1$)

$$\hat{H} = \Delta(\hat{n}_1 + \hat{n}_2) + U \sum_j \hat{a}_j^{\dagger 2} \hat{a}_j^2 + J(\hat{a}_1^{\dagger} \hat{a}_2 + \hat{a}_2^{\dagger} \hat{a}_1) + F_1(\hat{a}_1^{\dagger} + \hat{a}_1) + (F_2 \hat{a}_2^{\dagger} + F_2^* \hat{a}_2), \quad (1)$$

where Δ is the cavity-laser detuning and U the on-site Kerr nonlinearity. Each site decays at rate γ via Lindblad photon loss. We set $F_1 \in \mathbb{R}$, $F_2 = F_1 e^{i\phi}$, and $\tilde{E} \equiv \Delta - i\gamma/2$.

In the weak-drive limit $F_1 \ll \gamma$, quantum jump terms are suppressed by F_1/γ relative to the no-jump terms at each order in the Fock expansion (Appendix A), so the open-system dynamics reduces to a non-Hermitian Schrödinger equation. The steady state is therefore well approximated by a Fock-state amplitude expansion truncated at two photons [13], $|\psi_{ss}\rangle \approx |0, 0\rangle + \sum_{m+n \leq 2} C_{mn} |m, n\rangle$, with $g_{22}^{(2)}(0) = 2|C_{02}|^2/|C_{01}|^4$. UPB on site 2 requires $C_{02} = 0$.

III. ANALYTIC RESULTS

With $p \equiv e^{i\phi}$, the one-photon equations give $C_{01} = rC_{10}$ with $r = (J - p\tilde{E})/(pJ - \tilde{E})$ (see Appendix A). For $p = i$ ($\phi = 90^\circ$), $|r| < 1$ for all $J, \Delta, \gamma > 0$: site 2 is always less occupied than site 1. Imposing $C_{02} = 0$ in the two-photon sector, we derive the exact condition

$$J^2[2\tilde{E} + U(1 + p^2)] = 2p\tilde{E}(\tilde{E} + U)(2J - p\tilde{E}), \quad (2)$$

solved explicitly for U as

$$U = \frac{-2\tilde{E}(J - p\tilde{E})^2}{J^2(1 + p^2) - 4pJ\tilde{E} + 2p^2\tilde{E}^2}. \quad (3)$$

The physical solution (real, positive U) requires $\text{Im}(U) = 0$, determining $\Delta_{\text{opt}}(\phi, J)$. No solution exists at $\phi = 0$ or π : at these phases the two indirect pathways add constructively rather than destructively, precluding cancellation of C_{02} (Appendix B). For $J > \gamma/2$, a $U = 0$ solution appears at $\phi = \arcsin(\gamma/2J)$, but this is a *linear dark state* with $n_2 = 0$ identically—not genuine UPB (Appendix C).

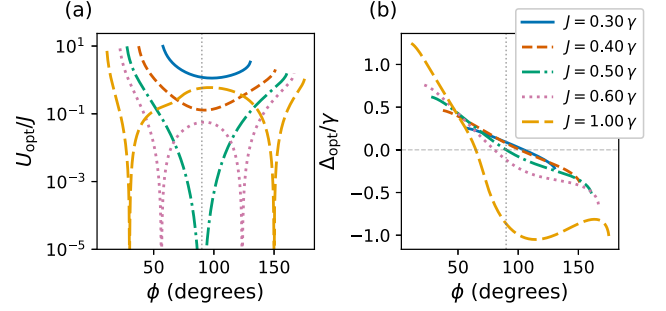


FIG. 2. UPB parameter locus vs drive phase ϕ . Left: required nonlinearity U_{opt}/J (log scale). Right: optimal detuning $\Delta_{\text{opt}}/\gamma$. There are four values of J/γ ; dotted vertical line, $\phi = 90^\circ$.

Setting $p = i$, the factor $1 + p^2$ vanishes and (2) collapses to the compact quadratic

$$(\tilde{E} + U)(\tilde{E} + 2iJ) = J^2, \quad (4)$$

whose real and imaginary parts give the closed-form parameter locus

$$\Delta_{\text{opt}} = \left(\frac{\gamma}{2} - J\right) \sqrt{\frac{4J}{\gamma} - 1}, \quad (5)$$

$$U_{\text{opt}} = \frac{4(J - \gamma/2)^2}{\sqrt{4J/\gamma} - 1}, \quad (6)$$

valid for $J > J_{\min} = \gamma/4$ —a factor of $2\sqrt{2} \approx 2.83$ below the single-site pumped UPB threshold $\gamma/\sqrt{2}$ [13].

The physical content of Eq. (4) is that the direct drive $F_2|0, 1\rangle \rightarrow |0, 2\rangle$ is exactly canceled by hopping from $|1, 1\rangle$ [Fig. 1(b)], regardless of how $|1, 1\rangle$ was populated. The Kerr shift U detunes the two-photon resonance, providing the degree of freedom needed for this cancellation to occur at a real detuning Δ ; the master condition (4) is the self-consistency requirement that the amplitude of $|1, 1\rangle$ is consistent with its equation of motion (Appendix A).

Solutions exist only in a range of ϕ that broadens with increasing J but is not symmetric around 90° , and U_{opt} is minimized close to $\phi = 90^\circ$ (within 0.7% of the global minimum for $J = 0.4\gamma$), making the quadrature drive the natural operating point when the Kerr interaction is most constrained (Fig. 2). Figure 3(a) shows how all three equal-time correlators vary with detuning Δ at the UPB optimum. $g_{11}^{(2)}(0)$ remains close to unity throughout, confirming that site 1 retains near-Poissonian statistics. $g_{22}^{(2)}(0)$ exhibits strong bunching for detunings below the optimum—reflecting two-photon resonance enhancement—before dropping sharply to zero at Δ_{opt} , where the interference condition is satisfied.

IV. OVERDAMPED REGIME AND THE TIME-DELAYED CORRELATOR

The threshold $J_{\min} = \gamma/4$ places the natural operating point at $J \lesssim \gamma/2$, the overdamped regime where γ exceeds the intermode oscillation frequency $2J$. The suppression of the rapid oscillations is therefore indirect: the quadrature drive does not itself modify the normal-mode dynamics, but lowers the UPB threshold from $\gamma/\sqrt{2}$ to $\gamma/4$, permitting operation

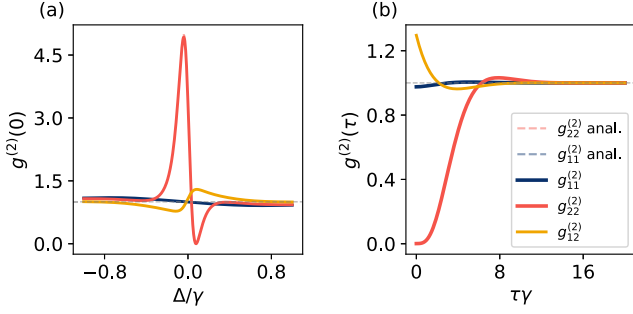


FIG. 3. Second-order correlators at the optimal point ($J = 0.4\gamma$, $U = 0.052\gamma$, $\Delta = 0.077\gamma$). Solid curves are full Lindblad master-equation simulations; dashed curves are the analytic results (indistinguishable on this scale). (a) Equal-time correlators $g_{11}^{(2)}(0)$ (navy), $g_{22}^{(2)}(0)$ (red), and $g_{12}^{(2)}(0)$ (yellow) vs detuning Δ at fixed $F_1 = 0.01\gamma$; dashed curves are the analytic results (indistinguishable from the numerics on this scale). (b) Time-delayed correlators $g_{11}^{(2)}(\tau)$ (navy), $g_{22}^{(2)}(\tau)$ (red), and $g_{12}^{(2)}(\tau)$ (yellow); dashed curves are the analytic results (indistinguishable from the numerics on this scale).

at $J < \gamma/2$, where the beat frequency $2J$ falls below the decay rate γ . We derive a closed-form expression for $g_{22}^{(2)}(\tau)$ via the quantum regression theorem in the weak-drive limit [20] (Appendix D). After detecting a photon from site 2, the postdetection state $\hat{a}_2|\psi_{ss}\rangle = C_{01}|0, 0\rangle + C_{11}|1, 0\rangle$ (using $C_{02} = 0$) lies entirely in the one-photon sector. It evolves under the driven one-photon equations, which diagonalize into two decoupled modes with decay rates $u_{1,2} = \gamma/2 \mp i\omega_{1,2}$ and oscillation frequencies $\omega_{1,2} = \Delta_{\text{opt}} \pm J$. The real part $\gamma/2$ is common to both normal modes and independent of J : the coherent coupling shifts only the mode frequencies (splitting them by $2J$) and introduces no additional loss. “Overdamped” here denotes $2J < \gamma$, i.e. the intermode beat is damped before completing an oscillation, in the sense of a damped harmonic oscillator. The result $g_{22}^{(2)}(\tau) = |C_{01}(\tau)|^2/|C_{01}|^4$ satisfies $g_{22}^{(2)}(0) = 0$ and $g_{22}^{(2)}(\infty) = 1$ exactly; the analytic $g_{11}^{(2)}(\tau)$ and $g_{12}^{(2)}(\tau)$ are shown in Fig. 3(b), where $g_{11}^{(2)}(\tau)$ remains close to unity at all delays while $g_{22}^{(2)}(\tau)$ rises smoothly from zero without rapid oscillations, crossing 0.5 at $\tau \approx 3.2/\gamma$.

The smooth rise of $g_{22}^{(2)}(\tau)$ to unity, free of rapid oscillations, is confirmed numerically. Quantitatively, by “oscillation-free” we mean that in the overdamped regime the transient consists of at most a single overshoot of unity, with no multi-peaked ringing at frequency $2J$; genuine oscillations appear only for $J > \gamma/2$. The peak overshoot $g_{\text{max}}^{(2)} - 1$ is nonmonotonic across J : it grows from 0.4% near threshold to 8% near $J = \gamma/2$ (3.3% at the operating point $J = 0.4\gamma$), then decreases in the underdamped regime, reaching 0.04% at $J = \gamma$ despite oscillation periods $T_2 \sim 3/\gamma$. The shrinking overshoot reflects decreasing transient amplitude relative to the steady-state value $|C_{01}|^2$ (Appendix F gives the full J -scan and tabulated values).

To see how the correlators behave at higher drive amplitudes and across a range of detunings, we solve the full master equation numerically (see Appendix E). At $F_1 = 0.01\gamma$ the simulations match the analytic results nearly perfectly (Fig. 3); the breakdown of the no-jump approximation with increasing F_1 is quantified in Appendix E (see also Fig. 7), where the minimum of $g_{22}^{(2)}(0)$ is seen to rise and redshift as F_1 increases. The full dependence on F_1 and Δ is shown in Fig. 4: $g_{11}^{(2)}(0)$ hovers close to unity throughout the landscape, while at the maximum drive $F_1 = 0.25\gamma$ the minimum $g_{22}^{(2)}(0)$ rises to ≈ 0.46 , with mean occupations $\langle n_1 \rangle \approx 0.30$ and $\langle n_2 \rangle \approx 7 \times 10^{-3}$.

V. PULSED EXCITATION

On-demand single-photon emission requires pulsed operation. Earlier UPB single-photon proposals operated under pulsed excitation either within the narrow $\delta\tau \sim 1/J$ antibunching window [16] or by removing the normal-mode oscillations through a nonreciprocal one-way dissipative coupling [17]. The overdamped regime accessed here achieves the same clean pulsed operation on a standard Hermitian dimer: the $\sim 1/\gamma$ antibunching window means the antibunching survives pulsed excitation without overshoots or oscillations. We simulate a hybrid drive: site 1 is pumped continuously at amplitude F_1 while site 2 receives a Gaussian pulse $F_2(t) = iF_1 e^{-t^2/2\sigma^2}$ with $F_1 = 0.05\gamma$. This models a cw pump that maintains a steady population in the dimer, with a pulsed

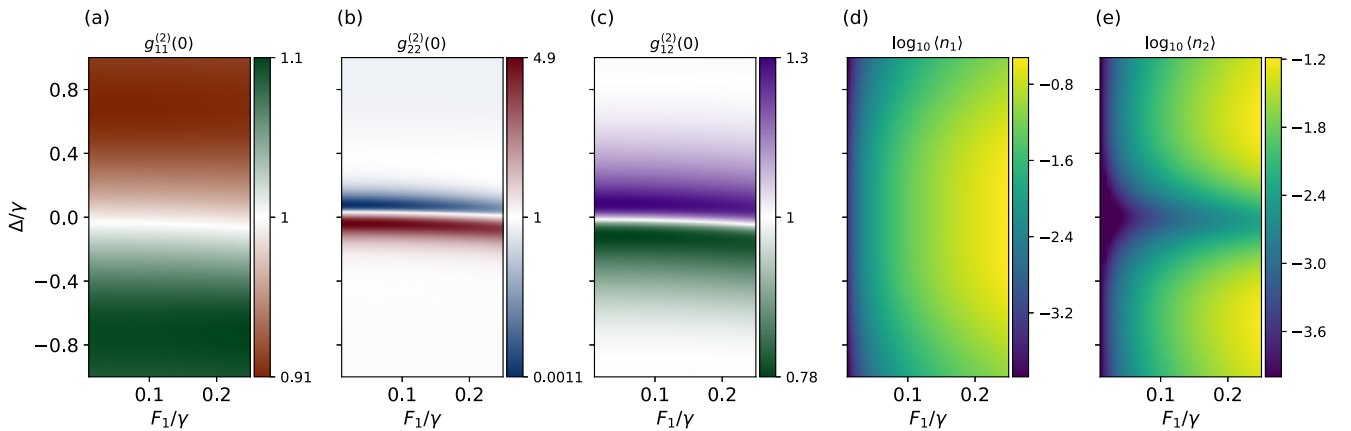


FIG. 4. Parameter landscape at $\phi = 90^\circ$, $J = 0.4\gamma$, $U = 0.052\gamma$, as a function of detuning Δ/γ (vertical) and drive amplitude $F_1/\gamma \in [0.01, 0.25]$ (horizontal). All panels are full Lindblad master-equation simulations. Panels from left: $g_{11}^{(2)}(0)$, $g_{22}^{(2)}(0)$, $g_{12}^{(2)}(0)$, $\log_{10}\langle n_1 \rangle$, $\log_{10}\langle n_2 \rangle$.

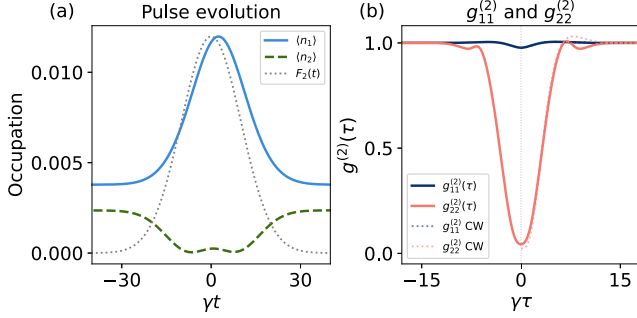


FIG. 5. Pulsed excitation ($J = 0.4\gamma$, $\sigma = 10/\gamma$, $F = 0.05\gamma$). Left: occupations $\langle n_1 \rangle$ (blue), $\langle n_2 \rangle$ (dashed green), and scaled pulse envelope (dotted). Right: $g_{11}^{(2)}(\tau)$ (navy) and $g_{22}^{(2)}(\tau)$ (red), with cw references (dotted). $\tau = 0$ is referenced to the pulse peak center ($t = 0$).

quadrature field triggering single-photon emission from site 2. For $\sigma = 10/\gamma$, the equal-time value $g_{22}^{(2)}(0)$ at the pulse peak approaches the cw dip, confirming that the antibunching is robust on the timescale of the pulse width and free of ringing (Fig. 5). For a GaAs quantum-well photonic molecule with $Q \simeq 10^4$ at $\lambda = 810$ nm, $1/\gamma \approx 4.3$ ps and the pulse width $\sigma = 10/\gamma \approx 43$ ps gives a repetition rate of ≈ 5 GHz. At $F_1 = 0.05\gamma$ the mean occupation reaches $\langle n_2 \rangle \approx 0.012$, giving a detected photon rate of ≈ 20 MHz with $g_{22}^{(2)}(0) < 0.05$ —competitive with state-of-the-art GaAs single-photon sources [8]. Because the $\sim 1/\gamma$ antibunching window spans essentially the whole emission, this rate requires no detector time gating; in conventional UPB the $\sim 1/J$ window forces gating onto a small fraction of the emission, reducing the usable rate accordingly.

VI. DISCUSSION

The bilateral quadrature drive achieves two decisive advantages over the single-site-driven UPB scheme [13]. First, the single-site scheme requires $J > \gamma/\sqrt{2}$ for any solution to exist; below this threshold no value of U can satisfy the interference condition. Our scheme operates for $J > \gamma/4$, opening the entire parameter space $\gamma/4 < J < \gamma/\sqrt{2}$ that is inaccessible to the single-site drive, including the overdamped regime $J < \gamma/2$ where $g^{(2)}(\tau)$ is smooth and oscillation free. At the operating point $J = 0.4\gamma$ we obtain $U \approx 0.052\gamma$ with no counterpart in the single-site scheme. For $J \gtrsim 0.96\gamma$ the single-site U eventually falls below ours (the two are equal at $J \approx 0.96\gamma$, $U \approx 0.50\gamma$), but that regime requires $J \gg \gamma$, placing extreme demands on cavity Q factors and yielding a rapidly oscillating $g^{(2)}(\tau)$ with antibunching window $\delta\tau \sim 1/J$. A minimal two-cavity geometry driven with a single laser and a 90° electro-optic phase shift is therefore sufficient to achieve UPB in the experimentally favorable overdamped regime.

The use of a controlled interport phase to enforce destructive interference in a two-mode system recalls coherent perfect absorption [25,26], where a tuned relative phase between two input ports suppresses the linear response; here the same principle is applied one level higher in the Fock hier-

archy, selectively canceling the two-photon amplitude while leaving the one-photon output intact.

The general condition (3) shows that $\phi = 90^\circ$ minimizes U_{opt} to within 0.7% of the global optimum for $J < \gamma/2$, making it the natural operating point when Kerr interaction strength is the limiting constraint. Along the locus in (5) and (6), U_{opt} has no interior minimum with $U > 0$: it decreases monotonically toward zero as $J \rightarrow (\gamma/2)^-$ (the linear dark-state boundary) and increases for $J > \gamma/2$. The operating point $J \approx 0.4\gamma$, giving $U_{\text{opt}} \approx 0.052\gamma$, balances a small required nonlinearity against sufficient distance from this boundary. Promising platforms include Cu₂O Rydberg exciton polaritons [27], semiconductor photonic molecules [28,29], trion polaritons in 2D monolayers [30], and moiré polariton systems [11] where enhanced nonlinearities are accessible.

A particularly compact implementation is possible within a *single* birefringent micropillar [31], where the H and V polarization modes play the roles of the two sites, the polarization splitting δ provides the hopping $J = \delta/2$, and the quadrature drive condition $F_2 = iF_1$ is satisfied by circularly polarized excitation—requiring only a quarter-wave plate in the beam path. The single-pillar geometry introduces an additional cross-Kerr interaction $U_x \hat{n}_H \hat{n}_V$ from opposite-spin exciton-exciton scattering, which one might worry could perturb the UPB condition. However, since U_x shifts only the $|1, 1\rangle$ two-photon state—leaving $|0, 2\rangle$ and $|2, 0\rangle$ unchanged—it enters the UPB condition only at first order through the self-consistency of C_{11} . For GaAs, where the opposite-spin interaction U_x is small and negative [32], U_{opt} shifts by less than 3% across the full experimentally reported range of U_x/U , while Δ_{opt} may shift by up to $\approx 10\%$ at the largest values, and the minimum coupling threshold $J_{\text{min}} = \gamma/4$ is unchanged to all orders in U_x/U . The single-pillar scheme is therefore as robust as the two-site dimer.

We quantify the sensitivity to fabrication disorder and its compensation in Appendix F: with fixed drive parameters the $g_{22}^{(2)}(0) < 0.1$ condition is maintained for detuning mismatches up to $\pm 0.033\gamma$ (≈ 1 GHz in GaAs), but retuning the drive phase ϕ alone extends this tolerance nearly tenfold to $\pm 0.26\gamma$, and adjusting both ϕ and the amplitude ratio $|F_2|/|F_1|$ restores exact UPB, eliminating the need for post-fabrication cavity trimming. Multiple identical dimers can be etched across the same semiconductor wafer in a single lithographic step, with low site-to-site disorder and near-identical emission frequencies, offering a route to parallelized single-photon generation with high photon indistinguishability that is difficult to achieve with self-assembled quantum dots without postfabrication tuning.

The dimer unit tiles naturally into a driven-dissipative lattice of symmetrically driven Kerr dimers, where the overdamped condition $J < \gamma/2$ relaxes quality-factor requirements compared to conventional UPB schemes, providing a practical starting point for exploring quantum many-body phases in weakly nonlinear photonic and polaritonic arrays.

In conclusion, equal-amplitude bilateral driving with a 90° phase difference reduces the UPB coupling threshold to $J_{\text{min}} = \gamma/4$ and places the system in the overdamped regime, yielding a smooth $g^{(2)}(\tau)$ directly resolvable with standard detectors. Unlike schemes requiring nonreciprocal couplings

or large photon networks, the approach uses only a single laser with a controlled phase shift and a standard Hermitian dimer, well suited to weakly nonlinear platforms such as Cu₂O Rydberg exciton polaritons [27] and semiconductor photonic molecules [28] where conventional photon blockade is beyond reach. Site-to-site fabrication disorder can be compensated by retuning the drive: adjusting the phase alone extends the disorder tolerance nearly tenfold, and adjusting both the phase and the amplitude ratio restores the blockade exactly, removing the need for postfabrication cavity trimming and opening a route to arrays of identical single-photon sources on a single chip. The minimal two-cavity geometry and overdamped operating condition make the scheme directly extensible to driven-dissipative lattices, providing a scalable platform for quantum many-body physics in weakly nonlinear photonic and polaritonic systems.

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DATA AVAILABILITY

The simulation scripts, numerical output data, and plotting code used to produce all figures in this manuscript are available in the University of St Andrews Research Portal [33].

APPENDIX A: DERIVATION OF THE $C_{02} = 0$ CONDITION AT $\phi = 90^\circ$

1. Hamiltonian and the Lindblad master equation

We consider two coupled single-mode Kerr cavities driven coherently in the rotating frame ($\hbar = 1$), with the Hamiltonian as in the main text [Eq. (1)]. Photon loss from each site at rate γ is modeled by the Lindblad master equation:

$$\dot{\rho} = -i[\hat{H}, \rho] + \gamma \left(\sum_{k=1,2} \hat{a}_k \rho \hat{a}_k^\dagger - \frac{1}{2} \{\hat{N}, \rho\} \right), \quad (\text{A1})$$

where $\hat{N} = \hat{n}_1 + \hat{n}_2$ is the total photon number operator, valid since both sites share the same decay rate γ . We use the *Kerr convention* $U \hat{a}^\dagger \hat{a}^2$ without an additional factor of $1/2$, matching Ref. [13], so the two-photon state $|2\rangle$ acquires an energy shift $2U$ relative to the noninteracting case.

2. No-jump limit and weak-drive Fock-state expansion

In the weak-drive limit $F_1 \ll \gamma$, density-matrix elements scale as $\rho_{mn,m'n'} \sim (F_1/\gamma)^{(m+n+m'+n')/2}$, since each photon costs one drive insertion suppressed by $1/\gamma$. The jump superoperator contributes to the equation of motion for $\rho_{mn,m'n'}$ via the source term

$$\gamma \sum_k \langle m, n | \hat{a}_k \rho \hat{a}_k^\dagger | m', n' \rangle \propto \rho_{m+1,n;m'+1,n'} + \rho_{m,n+1;m',n'+1}, \quad (\text{A2})$$

which samples ρ at total photon number $m+n+1$ (or $m'+n'+1$)—one order higher in F_1/γ than $\rho_{mn,m'n'}$ itself.

The jump contribution to the amplitude equations is therefore suppressed by F_1/γ relative to the no-jump terms and can be neglected throughout the truncated space. The open-system dynamics reduces to a non-Hermitian Schrödinger equation governed by the effective Hamiltonian

$$\hat{H}_{\text{eff}} = \hat{H} - i\frac{\gamma}{2}\hat{N}, \quad (\text{A3})$$

where $\hat{N} = \hat{n}_1 + \hat{n}_2$. This replaces the real cavity detuning Δ with the complex frequency $\tilde{E} \equiv \Delta - i\gamma/2$ throughout all amplitude equations, encoding the cavity decay as an imaginary part of the single-photon energy.

The steady state is therefore well approximated by the pure-state ansatz

$$|\psi_{\text{ss}}\rangle \simeq \sum_{n+m \leq 2} C_{nm} |n, m\rangle, \quad (\text{A4})$$

with $C_{00} \approx 1$. The correlators are $g_{11}^{(2)}(0) = 2|C_{20}|^2/|C_{10}|^4$ and $g_{22}^{(2)}(0) = 2|C_{02}|^2/|C_{01}|^4$; UPB on site 2 requires $C_{02} = 0$.

3. One-photon sector: Exact symmetry for $F_2 = iF_1$

With the amplitude equations governed by $\hat{H}_{\text{eff}}$ [Eq. (A3)], projecting onto $|1, 0\rangle$ and $|0, 1\rangle$ and setting $\tilde{C}_{10} = \tilde{C}_{01} = 0$ gives the linear system

$$\begin{pmatrix} \tilde{E} & J \\ J & \tilde{E} \end{pmatrix} \begin{pmatrix} C_{10} \\ C_{01} \end{pmatrix} = \begin{pmatrix} -F_1 \\ -F_2 \end{pmatrix}, \quad (\text{A5})$$

where $\tilde{E} = \Delta - i\gamma/2$ appears on both diagonal entries since both sites share the same detuning and decay rate. For $F_2 = iF_1$ one verifies immediately that $C_{01} = rC_{10}$ with

$$r = \frac{i\tilde{E} - J}{\tilde{E} - iJ}. \quad (\text{A6})$$

Since $(J + \gamma/2)^2 > (J - \gamma/2)^2$ for all $J, \gamma > 0$, we have $|r| < 1$ unconditionally: site 2 is always less occupied than site 1 despite $|F_2| = |F_1|$.

4. Two-photon sector and the UPB condition

Projecting onto $|2, 0\rangle$, $|1, 1\rangle$, and $|0, 2\rangle$ gives three coupled equations. Defining $E_U \equiv \tilde{E} + U$ and using the *crucial identity* $E_2 \equiv 2\Delta - i\gamma = 2\tilde{E}$ (valid for identical cavities), the steady-state equations are

$$2E_U C_{20} + \sqrt{2}J C_{11} = -\sqrt{2}F_1 C_{10}, \quad (\text{A7})$$

$$\sqrt{2}J C_{20} + 2\tilde{E} C_{11} + \sqrt{2}J C_{02} = -F_2 C_{10} - F_1 C_{01}, \quad (\text{A8})$$

$$\sqrt{2}J C_{11} + 2E_U C_{02} = -\sqrt{2}F_2 C_{01}. \quad (\text{A9})$$

Setting $C_{02} = 0$ in (A9) gives $C_{11} = -iF_1 C_{01}/J$. This has a direct physical meaning: the steady-state equation for C_{02} ,

$$2(\tilde{E} + U) C_{02} + \sqrt{2}J C_{11} = -\sqrt{2}F_2 C_{01}, \quad (\text{A10})$$

shows that $C_{02} = 0$ requires $J C_{11} = -F_2 C_{01}$: the direct drive $F_2|0, 1\rangle \rightarrow |0, 2\rangle$ is exactly canceled by hopping $J|1, 1\rangle \rightarrow |0, 2\rangle$. This cancellation fixes C_{11} regardless of how $|1, 1\rangle$ was reached. Substituting C_{11} into (A7) gives C_{20} ; inserting both into (A8) and imposing self-consistency of the one-photon

TABLE I. Parameter locus for the $\phi = 90^\circ$ bilateral-drive UPB. The $J = 0.30\gamma$ row has $U/J > 1$, which lies outside the spirit of weak-nonlinearity UPB but still satisfies the interference condition exactly.

J/γ	U_{opt}/γ	U_{opt}/J	$\Delta_{\text{opt}}/\gamma$
0.30	0.358	1.193	+0.089
0.40	0.052	0.129	+0.077
0.60	0.034	0.056	−0.118
0.70	0.119	0.170	−0.268
0.80	0.243	0.303	−0.445
1.00	0.577	0.577	−0.866

ratio r then yields the *main condition*:

$$(\tilde{E} + U)(\tilde{E} + 2iJ) = J^2. \quad (\text{A11})$$

The factor $2iJ$ has two origins: the $\sqrt{2}J$ two-photon hopping matrix elements (giving the factor of 2) and the 90° drive phase $F_2 = iF_1$ (giving the factor of i), which rotates the effective coupling into the imaginary direction. The simplification from the general quartic of Ref. [20] to a quadratic relies on $E_2 = 2\tilde{E}$, which holds only for identical cavities.

5. Solving for the optimal parameter locus

Writing $\tilde{E} = \Delta - i\gamma/2$ with Δ real and imposing $\text{Im}[E_U(\tilde{E} + 2iJ)] = 0$ (since the right-hand side J^2 is real), one obtains

$$U = \frac{2\Delta(\gamma - 2J)}{4J - \gamma}. \quad (\text{A12})$$

Substituting into the real part of (A11) and solving for Δ gives the closed-form locus [Eqs. (5) and (6) of the main text]:

$$\Delta_{\text{opt}} = \left(\frac{\gamma}{2} - J\right) \sqrt{\frac{4J}{\gamma} - 1}, \quad (\text{A13})$$

$$U_{\text{opt}} = \frac{4(J - \gamma/2)^2}{\sqrt{4J/\gamma} - 1}, \quad (\text{A14})$$

valid for $J > J_{\text{min}} = \gamma/4$. Selected numerical values are given in Table I.

APPENDIX B: GENERAL DRIVE PHASE: DERIVATION AND PHASE RANGE

1. General $C_{02} = 0$ condition for arbitrary ϕ

For $F_2 = F_1 e^{i\phi}$ with $p \equiv e^{i\phi}$, the one-photon amplitudes are

$$C_{10} = \frac{F_1(pJ - \tilde{E})}{\tilde{E}^2 - J^2}, \quad C_{01} = \frac{F_1(J - p\tilde{E})}{\tilde{E}^2 - J^2}, \quad (\text{B1})$$

with ratio $r = (J - p\tilde{E})/(pJ - \tilde{E})$. Following the same two-photon projection procedure as Appendix A and using $E_2 = 2\tilde{E}$, one derives

$$J^2[2\tilde{E} + U(1 + p^2)] = 2p\tilde{E} E_U(2J - p\tilde{E}), \quad (\text{B2})$$

solved explicitly for U as

$$U = \frac{-2\tilde{E}(J - p\tilde{E})^2}{J^2(1 + p^2) - 4pJ\tilde{E} + 2p^2\tilde{E}^2}. \quad (\text{B3})$$

TABLE II. UPB operating points as a function of drive phase for $J = 0.4\gamma$ and values from Eq. (B3) with $\text{Im}(U) = 0$.

ϕ	U_{opt}/γ	$\Delta_{\text{opt}}/\gamma$
40°	1.835	+0.453
60°	0.192	+0.317
90°	0.052	+0.077
120°	0.113	−0.112
150°	0.663	−0.351

Setting $p = i$ makes $1 + p^2 = 0$ and the denominator reduces to $-2\tilde{E}(\tilde{E} + 2iJ)$, recovering (A11) exactly.

2. Failure at $\phi = 0$ and $\phi = \pi$

At $p = 1$ ($\phi = 0$), Eq. (B3) gives $U = -\tilde{E}$, which requires $\text{Im}(U) = \gamma/2 \neq 0$ —no physical solution exists. The same holds at $p = -1$ ($\phi = \pi$). Destructive interference in the two-photon sector demands a nonzero quadrature component; a purely in-phase or antiphase drive cannot produce $C_{02} = 0$.

3. Valid phase range and the minimum of U_{opt}

Physical solutions (U real and positive) exist only in a finite range of ϕ that is not symmetric about 90° . For $J = 0.4\gamma$ this range is approximately 38° to 151° .

The minimum of $U_{\text{opt}}(\phi)$ at fixed J occurs at $\phi^* \approx 92.2^\circ$ for $J = 0.4\gamma$, only 0.7% below the value at $\phi = 90^\circ$. The phase ϕ^* approaches 90° as $J \rightarrow (\gamma/2)^-$. Table II lists U_{opt}/γ and $\Delta_{\text{opt}}/\gamma$ for representative phases at $J = 0.4\gamma$.

APPENDIX C: LINEAR DARK STATE

For $J > \gamma/2$, Eq. (B3) admits a spurious solution with $U = 0$ at the phase $\sin \phi = \gamma/(2J)$. At this point $J = p\tilde{E}$, which causes $C_{01} = 0$ exactly from (B1): the direct drive on site 2 is canceled by the hopping amplitude, leaving site 2 unpopulated at the *single-photon* level. Consequently $C_{02} = 0$ follows trivially—there is no photon population on site 2 to build on—and $g_{22}^{(2)}(0) = 0/0$ is undefined rather than a quantum antibunching signature.

APPENDIX D: ANALYTIC $g_{22}^{(2)}(\tau)$ VIA THE QUANTUM REGRESSION THEOREM

1. Postdetection state

At the UPB operating point $C_{02} = 0$, the state immediately after detecting a photon from site 2 is, in the weak-drive limit,

$$|\tilde{\Psi}(0)\rangle = C_{01}|0, 0\rangle + C_{11}|1, 0\rangle, \quad (\text{D1})$$

which lies entirely in the zero- and one-photon sectors. Hence $g_{22}^{(2)}(0) = 0$ exactly at the UPB point.

2. Equations of motion and diagonalization

Under the cw bilateral drive, the one-photon amplitudes $C_{10}(\tau)$ and $C_{01}(\tau)$ obey the driven linear system (A5) (in time-dependent form with the vacuum amplitude fixed at C_{01} following Ref. [20]). Defining sum and difference variables $x = C_{10} + C_{01}$ and $y = C_{10} - C_{01}$, these decouple into two

first-order ordinary differential equations with decay rates and oscillation frequencies:

$$u_{1,2} = i(\tilde{E} \pm J) = \frac{\gamma}{2} \mp i(\Delta \pm J). \quad (\text{D2})$$

Both modes decay since $\text{Re}(u_{1,2}) = \gamma/2 > 0$. In the overdamped regime $J < \gamma/2$, the imaginary parts of $u_{1,2}$ have opposite signs and the two exponentials interfere, producing a transient with at most a single small overshoot of unity and no oscillations.

3. Closed-form result

With initial conditions $C_{10}(0) = C_{11}$, $C_{01}(0) = 0$, the solution is

$$C_{01}(\tau) = \frac{1}{2} \left[\left(C_{11} - \frac{\lambda_1}{u_1} \right) e^{-u_1 \tau} - \left(C_{11} - \frac{\lambda_2}{u_2} \right) e^{-u_2 \tau} + \frac{\lambda_1}{u_1} - \frac{\lambda_2}{u_2} \right], \quad (\text{D3})$$

where $\lambda_{1,2} = -i(F_1 \pm F_2)C_{01}$. The final result $g_{22}^{(2)}(\tau) = |C_{01}(\tau)|^2 / |C_{01}|^4$ gives $g_{22}^{(2)}(0) = 0$ and $g_{22}^{(2)}(\infty) = 1$ exactly, as verified by substituting the steady-state relations ($\tilde{E} \pm J)(C_{10} \pm C_{01}) = -(F_1 \pm F_2)$.

4. Oscillation frequencies at the UPB locus

Substituting Δ_{opt} from (A13), the oscillation frequencies become

$$\omega_{1,2} = \left(\frac{\gamma}{2} - J \right) \sqrt{\frac{4J}{\gamma} - 1 \pm J}. \quad (\text{D4})$$

For $J < \gamma/2$, the arguments of both exponentials in (D3) have positive real parts, confirming overdamped, oscillation-free behavior.

5. Analytic $g_{11}^{(2)}(\tau)$

For completeness, $g_{11}^{(2)}(\tau)$ is obtained analogously by applying the quantum regression theorem after a site-1 detection. The postdetection state is $C_{10}|0, 0\rangle + \sqrt{2}C_{20}|1, 0\rangle + C_{11}|0, 1\rangle$, giving initial conditions $C_{10}(0) = \sqrt{2}C_{20}$, $C_{01}(0) = C_{11}$. The result is $g_{11}^{(2)}(\tau) = |C_{10}(\tau)|^2 / |C_{10}|^4$, where

$$C_{10}(\tau) = \frac{1}{2} \left[\left(x_0 - \frac{\lambda_1^{(11)}}{u_1} \right) e^{-u_1 \tau} + \left(y_0 - \frac{\lambda_2^{(11)}}{u_2} \right) e^{-u_2 \tau} + \frac{\lambda_1^{(11)}}{u_1} + \frac{\lambda_2^{(11)}}{u_2} \right], \quad (\text{D5})$$

with $\lambda_{1,2}^{(11)} = -i(F_1 \pm F_2)C_{10}$, $x_0 = \sqrt{2}C_{20} + C_{11}$, $y_0 = \sqrt{2}C_{20} - C_{11}$. At our $C_{02} = 0$ operating point, $C_{20} \neq 0$. From Eq. (A7) with $C_{11} = -F_2 C_{01}/J$ and $F_2 = iF_1$, $C_{01} = rC_{10}$,

$$C_{20} = \frac{F_1 C_{10}(ir - 1)}{\sqrt{2}(\tilde{E} + U)}. \quad (\text{D6})$$

Substituting into $g_{11}^{(2)}(0) = 2|C_{20}|^2 / |C_{10}|^4$ using the initial condition $C_{10}(0) = \sqrt{2}C_{20}$,

$$g_{11}^{(2)}(0) = \frac{2|C_{20}|^2}{|C_{10}|^4} = \frac{|ir - 1|^2 F_1^2}{|\tilde{E} + U|^2 |C_{10}|^2}. \quad (\text{D7})$$

Numerically at the UPB optimum ($J = 0.4\gamma$), this gives $g_{11}^{(2)}(0) \approx 0.98$, confirming that antibunching is selective to site 2 while site 1 remains near Poissonian.

APPENDIX E: NUMERICAL METHODS

All numerical results were obtained using QUTIP 4 (Quantum Toolbox in PYTHON) [34] in PYTHON 3. The Hilbert space was truncated at $N_{\text{cut}} = 7$ photons per site (total dimension 49), verified to be sufficient by checking convergence at $N_{\text{cut}} = 15$.

1. Steady-state correlators

Equal-time correlators $g_{jk}^{(2)}(0)$ and mean occupations $\langle n_j \rangle$ were computed from the steady-state density matrix ρ_{ss} obtained via `qutip.steadystate` using the direct sparse LU solver. For the parameter landscape (Fig. 4) the grid was evaluated in parallel across all CPU cores using PYTHON's multiprocessing module. In the weak-drive limit $F_1 \ll \gamma$, the perturbative amplitude solver (Appendix A) was used instead for speed, with the full master equation reserved for verification.

2. Time-delayed correlators

$g_{jk}^{(2)}(\tau)$ was computed via the quantum regression theorem (QRT) implemented directly with `qutip.mesolve`. For each correlator, the postmeasurement state $\hat{a}_k \rho_{\text{ss}} \hat{a}_k^\dagger$ was normalized and propagated forward in τ under the full Lindblad master equation; the expectation value of $\hat{a}_j^\dagger \hat{a}_j$ on the propagated state then yields $g_{jk}^{(2)}(\tau)$.

3. Pulsed excitation

The time-dependent master equation with a Gaussian drive envelope was integrated using the variable-step Adams-Bashforth-Moulton solver in `qutip.mesolve`. $g_{22}^{(2)}(\tau)$ was then computed via the QRT with $\tau = 0$ referenced to the pulse peak.

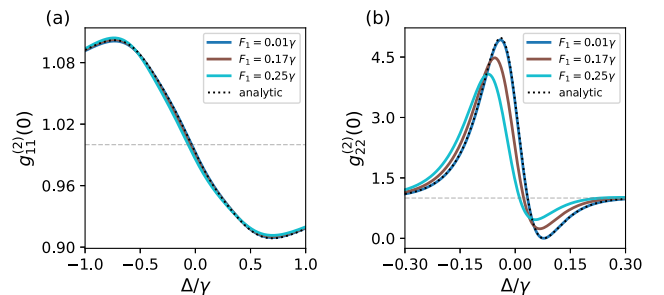


FIG. 6. Equal-time correlators $g_{11}^{(2)}(0)$ (left) and $g_{22}^{(2)}(0)$ (right) vs detuning Δ/γ at $J = 0.4\gamma$, $U = 0.052\gamma$, $\phi = 90^\circ$, for three drive amplitudes $F_1/\gamma \in \{0.01, 0.17, 0.25\}$ (solid coloured curves, full master equation). The black dotted curve is the analytic no-jump result, which is independent of F_1 in the weak-drive limit.

TABLE III. Oscillation frequencies and peak overshoot of $g_{22}^{(2)}(\tau)$ at the $\phi = 90^\circ$ UPB optimum. $T_{1,2} = 2\pi/|\omega_{1,2}|$ are the oscillation periods. The row at $J = \gamma/2$ is the transition between overdamped and underdamped regimes; $U_{\text{opt}} = 0$ at this point corresponds to the linear dark state and is excluded in practice.

J/γ	$\Delta_{\text{opt}}/\gamma$	U_{opt}/γ	ω_1/γ	ω_2/γ	$T_2\gamma$	$g_{\text{max}}^{(2)}$
0.26	+0.048	1.152	+0.308	-0.212	29.7	1.004
0.30	+0.089	0.358	+0.389	-0.211	29.8	1.006
0.40	+0.077	0.052	+0.478	-0.323	19.5	1.033
0.55	-0.055	0.009	+0.495	-0.605	10.4	1.113
0.60	-0.118	0.034	+0.482	-0.718	8.7	1.124
0.70	-0.268	0.119	+0.432	-0.968	6.5	1.087
0.80	-0.445	0.243	+0.355	-1.245	5.0	1.037
0.90	-0.645	0.397	+0.255	-1.545	4.1	1.014
1.00	-0.866	0.577	+0.134	-1.866	3.4	1.000

4. Validity of the no-jump approximation

Figure 6 compares the full master-equation result for $g_{11}^{(2)}(0)$ and $g_{22}^{(2)}(0)$ as a function of detuning Δ at three drive amplitudes ($F_1/\gamma = 0.01, 0.17, 0.25$) with the analytic no-jump prediction (black dotted curve), which is independent of F_1 in the weak-drive limit. At $F_1 = 0.01\gamma$ the master-equation result is indistinguishable from the analytic curve. As F_1 increases, the minimum of $g_{22}^{(2)}(0)$ rises and redshifts, reflecting corrections of order $(F_1/\gamma)^2$ beyond the no-jump approximation.

APPENDIX F: PEAK OVERSHOOT OF $g_{22}^{(2)}(\tau)$ ACROSS ALL J

Using the analytic formula (D3), Table III gives the oscillation frequencies and numerically computed peak value $g_{\text{max}}^{(2)} \equiv \max_{\tau>0} g_{22}^{(2)}(\tau)$ across the full valid range $J > \gamma/4$.

The peak overshoot $g_{\text{max}}^{(2)} - 1$ is nonmonotonic. In the overdamped regime ($J < \gamma/2$) it grows from $\approx 0.4\%$ near threshold to $\approx 8\%$ approaching $J = \gamma/2$. In the underdamped regime ($J > \gamma/2$) it decreases monotonically, reaching 0.04% at $J = \gamma$ despite oscillation periods $T_2 \sim 3/\gamma$. The shrinking overshoot reflects decreasing transient amplitudes

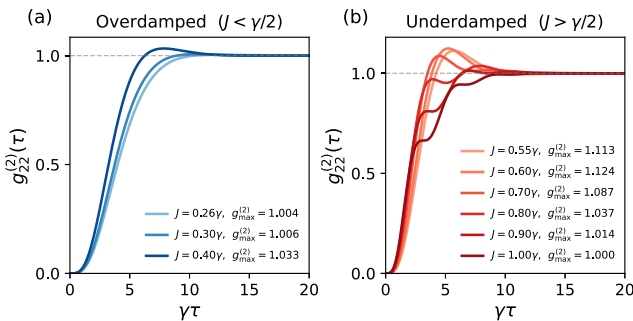


FIG. 7. Analytic $g_{22}^{(2)}(\tau)$ at the UPB optimum ($\phi = 90^\circ$, $F_1 = 0.05\gamma$) for various J/γ . (a) Blue panels ($J < \gamma/2$): smooth overdamped rise, overshooting unity by at most 8% . (b) Red panels ($J > \gamma/2$): visible oscillations at frequencies $\omega_{1,2} = \Delta_{\text{opt}} \pm \gamma$, with peak overshoot decreasing monotonically to 0.04% at $J = \gamma$.

TABLE IV. Disorder tolerances at the UPB optimum without compensation.

Disorder axis	Tolerance
Detuning $\delta\Delta$	$\pm 0.033\gamma$
Loss rate $\delta\gamma$	$\pm 0.060\gamma$
Nonlinearity δU	$\pm 0.033\gamma$

($C_{11} - \lambda/u$) in (D3) relative to the steady-state $|C_{01}|^2$ as J grows.

Figure 7 shows the analytic $g_{22}^{(2)}(\tau)$ across the full valid range, illustrating the crossover from overdamped (smooth) to underdamped (oscillatory) behavior.

APPENDIX G: DISORDER ROBUSTNESS

We assess the sensitivity of $g_{22}^{(2)}(0)$ to site-to-site parameter mismatch using the perturbative amplitude solver extended to asymmetric cavities. Three disorder axes are explored independently at the analytic optimum ($J = 0.4\gamma$, $U = 0.05164\gamma$, $\Delta = 0.07746\gamma$, $\phi = 90^\circ$, $F_1 = 0.05\gamma$): detuning mismatch $\delta\Delta = \Delta_1 - \Delta_2$, loss-rate mismatch $\delta\gamma = \gamma_1 - \gamma_2$, and nonlinearity mismatch $\delta U = U_1 - U_2$. We define the tolerance as the mismatch at which $g_{22}^{(2)}(0)$ exceeds 0.1 .

The tightest constraint is the detuning tolerance $\pm 0.033\gamma$ (Table IV). For a GaAs photonic molecule with $Q \simeq 10^5$ at $\lambda = 850$ nm ($\gamma \approx 25$ GHz), this corresponds to a frequency mismatch of ≈ 0.8 GHz (≈ 3 μeV). The loss-rate mismatch tolerance is roughly twice as large.

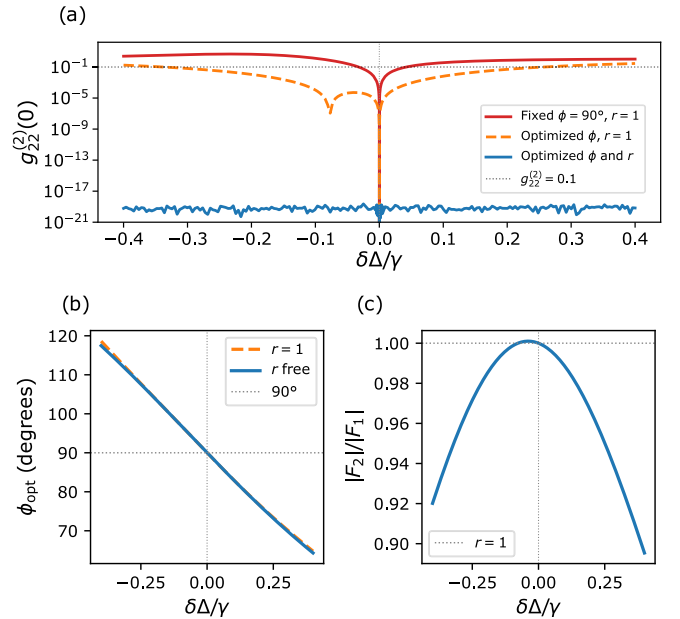


FIG. 8. Drive-parameter compensation of detuning disorder ($J = 0.4\gamma$, $U = 0.05164\gamma$, $\Delta = 0.07746\gamma$). (a) $g_{22}^{(2)}(0)$ vs $\delta\Delta$ for three cases: fixed $\phi = 90^\circ$ and $r = 1$ (red), optimized ϕ with $r = 1$ (orange dashed), and optimized ϕ and r (blue). Dotted line: $g_{22}^{(2)} = 0.1$ threshold. (b) Optimal ϕ as a function of $\delta\Delta$ for both compensated cases. (c) Optimal amplitude ratio $r = |F_2|/|F_1|$ for full compensation.

Compensation of disorder by drive retuning

The tolerances above assume fixed drive parameters. Since the $C_{02} = 0$ condition is one complex equation in three real system parameters, adding site-to-site disorder introduces new unknowns that can in principle be absorbed by retuning the drive. We investigate two strategies: (i) adjusting the drive phase ϕ alone while keeping $|F_2| = |F_1|$, and (ii) adjusting both ϕ and the amplitude ratio $r = |F_2|/|F_1|$.

For each value of detuning mismatch $\delta\Delta$, we minimize $g_{22}^{(2)}(0)$ numerically over the drive parameters. Figure 8 shows the result. Optimizing ϕ alone extends the tolerance from $\pm 0.033\gamma$ to $\pm 0.26\gamma$ —nearly an order of magnitude—

because ϕ directly controls the relative phase between the direct and indirect two-photon pathways. Optimizing both ϕ and r achieves $g_{22}^{(2)}(0) \lesssim 10^{-18}$ across the entire scanned range $|\delta\Delta| \leq 0.4\gamma$, demonstrating that the UPB condition can be exactly restored for arbitrary detuning disorder by drive retuning alone, with no need for postfabrication cavity trimming.

Figure 8(b) shows that the compensating ϕ tracks the disorder nearly linearly, deviating from 90° at a rate of approximately 15° per 0.1γ of mismatch. When r is also free, it decreases smoothly from 1 as the mismatch grows [Fig. 8(c)], indicating that the two pathways require unequal drive amplitudes to maintain destructive interference in the asymmetric dimer.

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