

Floquet engineering magnetism and superconductivity in the square-lattice Hubbard modelJan-Niklas Herre^{1,*}, Takuya Okugawa^{2,3}, Ammon Fischer³, Christoph Karrasch⁴, and Dante M. Kennes^{1,3}¹*Institute for Theory of Statistical Physics, RWTH Aachen University, Aachen, Germany*²*Department of Physics, Columbia University, New York, New York 10027, USA*³*Max Planck Institute for the Structure and Dynamics of Matter, Center for Free Electron Laser Science, Luruper Chaussee 149, 22761 Hamburg, Germany*⁴*Technische Universität Braunschweig, Institut für Mathematische Physik, Mendelssohnstraße 3, 38106 Braunschweig, Germany*

(Received 21 November 2025; revised 27 January 2026; accepted 17 February 2026; published 13 March 2026)

We study the interplay of magnetic order and superconductivity in the square-lattice Hubbard model under periodic driving with circularly polarized light. Formulating diagrammatic techniques based on the random-phase approximation in terms of Floquet Green's functions, allows us to analyze fluctuation-driven unconventional pairing for weak to moderate interactions. The interplay of repulsive interactions and photoassisted hopping of electrons gives rise to a rich magnetic phase diagram featuring an antiferromagnetic-to-ferromagnetic transition prior to a Floquet Lifshitz transition. Close to the antiferromagnetic transition, topological $d + id$ -wave superconductivity prevails in the phase diagram for a wide range of drive parameters. At intermediate- to high-frequency driving near the Floquet Lifshitz transition, superconducting orders are tuned from spin-singlet d -wave to spin-triplet p -wave character, providing an effective protocol for Floquet engineering topological superconductivity.

DOI: [10.1103/7cj5-tqys](https://doi.org/10.1103/7cj5-tqys)

The ability to manipulate quantum materials using tailored laser fields has opened new directions for dynamically controlling their properties [1]. In particular, Floquet engineering—the use of periodic driving to modify the effective low-energy Hamiltonian—has emerged as a promising route to realize exotic nonequilibrium states, ranging from topological phases [2–9] and magnetic order [10–13] to light-induced superconductivity [14–33]. Although many theoretical approaches focus on the high-frequency or low-amplitude limit, where heating is minimal and analytic control is tractable, the low- to intermediate-frequency regime at arbitrary drive amplitude remains less understood. Recent experiments further demonstrate that Floquet engineering of correlated materials extends beyond Mott insulators [8], underscoring the need for theoretical methods capable of addressing the weak- to moderate-coupling regime under subbandwidth driving, with direct relevance to Moiré-based systems [9].

We explore the periodically driven square-lattice Hubbard model, where previous works have investigated the possibility of Floquet engineering magnetic order in the strongly interacting limit [10,13] and, for weak coupling, suppression of antiferromagnetic (AFM) order at low drive frequencies [34]. In the limit of strong interactions, there has also been

progress on the Floquet engineering of chiral d -wave superconductivity [26,35]. Here, we complement these results using the nonequilibrium Floquet Green's function formalism [36–40] that can be systematically extended into the low- to intermediate-frequency regime beyond the reach of analytical expansions. To assess the emergence of collective charge, spin, and superconducting fluctuations, we resort to diagrammatic expansions on the level of the random-phase approximation. This enables us to study the impact of driven fermiology on the formation of electronic order. We find that, in the t - t' Hubbard model away from half filling driven by circularly polarized light at subbandwidth drive frequency, photoassisted processes lead to a switching between AFM and ferromagnetic (FM) orders. This is fundamentally different from the Floquet Lifshitz transition effects expected in the high-frequency limit, where next-to-nearest-neighbor hopping can be tuned to dominate over nearest-neighbor hopping. Close to the Floquet engineered magnetic instabilities, we find significant tunability of chiral d/p -wave superconductivity throughout the subbandwidth driving regime.

Model. We consider the square-lattice Hubbard model driven by an in-plane electric field and coupled to a metallic reservoir to facilitate a nonequilibrium steady state. The time-dependent Hamiltonian is given by

$$H(t) = - \sum_{ij,\sigma} \tilde{t}_{ij} e^{-i\mathbf{A}(t) \cdot \mathbf{R}_{ij}} \hat{c}_{i\sigma}^\dagger \hat{c}_{j\sigma} + U \sum_i \hat{n}_{i\uparrow} \hat{n}_{i\downarrow} + H_{\text{res}}, \quad (1)$$

where $\hat{c}_{i\sigma}^\dagger$ ($\hat{c}_{i\sigma}$) is the creation (annihilation) operator of a fermion with spin σ on site i and $\hat{n}_{i\sigma} = \hat{c}_{i\sigma}^\dagger \hat{c}_{i\sigma}$. U and $\tilde{t}_{ij}$ are the on-site repulsion and hopping between site i and j , respectively. For the latter, we allow for nearest-neighbor

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hopping $\tilde{t}$ as well as next-nearest-neighbor hopping $\tilde{t}'$ to break the perfect nesting condition and allow for superconducting order close to the magnetic transition [41]. Circularly polarized light, breaking time-reversal symmetry, is included via Peierls substitution [42] with its vector potential

$$\mathbf{A}(t) = \frac{E_0}{\Omega}(\cos \Omega t, \sin \Omega t), \quad (2)$$

where E_0 is the electrical field strength and Ω the drive frequency (drive period $T = 2\pi/\Omega$). The reservoirs are modeled by noninteracting tight-binding chains that are weakly coupled via a constant hopping term V to each lattice site independently,

$$H_{\text{res}} = \sum_{\mathbf{k}_r} \varepsilon_{\mathbf{k}_r} \hat{d}_{\mathbf{k}_r}^\dagger \hat{d}_{\mathbf{k}_r} + \sum_{\mathbf{k}, \mathbf{k}_r, \sigma} V (\hat{c}_{\mathbf{k}\sigma}^\dagger \hat{d}_{\mathbf{k}_r} + \hat{d}_{\mathbf{k}_r}^\dagger \hat{c}_{\mathbf{k}\sigma}). \quad (3)$$

In the wideband limit [43], this is modeled as a constant hybridization $\gamma = \pi \rho_0 |V|^2$ with ρ_0 being the reservoir density of states. The (thermal) reservoir is defined by the reservoir distribution function $f(\omega) = [e^{(\omega - \mu)/T_{\text{res}}} + 1]^{-1}$ at temperature T_{res} and chemical potential μ . The hybridization enters as an imaginary self-energy through the Dyson equation into the reservoir dressed Green's function of the system.

The system is initially in the uncorrelated normal state (paramagnetic), and we investigate how driving affects collective fluctuations and the emergence of magnetic and superconducting order in the weak to moderately interacting regime. The weak coupling to the bath facilitates the formation of a Floquet steady state that is synchronized with the drive (on the timescale $t^* \sim \gamma^{-1}$) [44]. In this formalism, dissipation and heating balance each other and no other heating channels are available.

High-frequency limit. In the limit of high driving frequency we can gain analytical insight via the Magnus expansion [45], which expands in $1/\Omega$ to find an effectively static Floquet Hamiltonian. The lowest-order result for the Hubbard model is known [15, 16, 46, 47], and we can define effective hopping parameters $\tilde{t}_{\text{eff}} \approx \tilde{t} \mathcal{J}_0(E_0/\Omega)$ and $\tilde{t}'_{\text{eff}} \approx \tilde{t}' \mathcal{J}_0(\sqrt{2}E_0/\Omega)$. In this limit, we can find two characteristic points. First, a Floquet Lifshitz transition [48–51] at $E_0/\Omega \approx 2.28$, where $\tilde{t}_{\text{eff}} = \tilde{t}'_{\text{eff}}$, which leads to a topological change in the Fermi surface from a continuous one to separated pockets. Second, the dynamical localization point lies close to the Floquet Lifshitz transition [52], where $\tilde{t}_{\text{eff}} \rightarrow 0$, at $E_0/\Omega \approx 2.4$. Here, the effective model is a square-lattice model on an enlarged unit cell with hopping $\tilde{t}'_{\text{eff}}$ and tilted by $\pi/4$.

Floquet random-phase approximation. Based on the Floquet Green's function formalism [36–40, 53, 54], we start from the retarded noninteracting $\text{SU}(2)$ -symmetric Floquet particle-hole susceptibility

$$\chi_{nn'}^{\text{ph}}(Q) = \frac{i}{N} \sum_{l, \mathbf{k}} \int \frac{d\omega}{2\pi} G_{n, n'+l}^{\text{R}}(K) G_{l, 0}^<(K - Q) + G_{n, n'+l}^<(K) G_{l, 0}^{\text{A}}(K - Q), \quad (4)$$

introducing the multi-index $K = (\omega, \mathbf{k})$ for the real frequency ω and the momentum $\mathbf{k}$ as well as discrete Floquet indices n, n', l . The Keldysh Green's functions that fulfill $G_{\mathbf{k}}^X(t, t') = G_{\mathbf{k}}^X(t + T, t' + T)$ for $X \in \{\text{R}, \text{A}, <\}$ were Fourier

transformed twice into

$$G_{K, nn'}^X = \frac{1}{T} \int_0^T d\bar{t} \int_{-\infty}^{\infty} d\tau e^{i(\omega + n\Omega)\bar{t} - i(\omega + n'\Omega)\tau'} G_{\mathbf{k}}^X(t, t'), \quad (5)$$

with $\bar{t} = \frac{t+t'}{2}$ and $\tau = t - t'$ [see Supplemental Material (SM) for details [55]]. Equation (4) is a matrix in Floquet space $\hat{\chi}_Q^{\text{ph}}$ and we can define the spin (s) and charge (c) susceptibility as a geometric sum within the random-phase approximation

$$\hat{\chi}^{s/c}(Q) = [\mathbb{1} \mp \hat{V}_0 \hat{\chi}^{\text{ph}}(Q)]^{-1} \hat{\chi}^{\text{ph}}(Q), \quad (6)$$

where the bare (and undriven) interaction is introduced as $\hat{V}_0 = U \mathbb{1}$ with $\mathbb{1}$ the identity matrix. The geometric series of the spin susceptibility diverges at a critical interaction strength $U_c = 1/\lambda_m$ where λ_m is the eigenvalue with maximal real part of all matrices $\hat{\chi}^{\text{ph}}(\mathbf{q}, \omega = 0)$, which is the Stoner criterion for the onset of magnetic order [56]. As long as $U < U_c$, the system remains paramagnetic and subsidiary superconducting order driven by spin-fluctuation exchange [57–61] can arise, captured by the singlet (s) and triplet (t) components of the retarded effective pairing vertex

$$\hat{V}^{\text{P},(s)}(Q) = \hat{V}_0 + \frac{3}{2} \hat{V}_0^2 \hat{\chi}^s(Q) - \frac{1}{2} \hat{V}_0^2 \hat{\chi}^c(Q), \quad (7)$$

$$\hat{V}^{\text{P},(t)}(Q) = -\frac{1}{2} \hat{V}_0^2 \hat{\chi}^s(Q) - \frac{1}{2} \hat{V}_0^2 \hat{\chi}^c(Q). \quad (8)$$

Unconventional superconducting order facilitated by nearby magnetic order is analyzed by diagonalizing a linearized gap equation using a static approximation of the vertex $\hat{V}_{\mathbf{k}-\mathbf{k}'}^{\text{P},(s)/(t)} \equiv \hat{V}^{\text{P},(s)/(t)}(\mathbf{k} - \mathbf{k}', \omega = 0)$. The superconducting gap equation in the Floquet-Keldysh formalism is given by

$$\Delta_{\mathbf{k}; n}^{(s)/(t)} = \frac{i}{N} \sum_{\mathbf{k}'; l} V_{\mathbf{k}-\mathbf{k}'; n-l}^{\text{P},(s)/(t)} \int \frac{d\omega'}{2\pi} \mathcal{F}_{\mathbf{k}'; l}^<(\omega'). \quad (9)$$

Here, we transformed the Floquet matrices into the single index notation from Ref. [36]. In linear approximation, the anomalous Green's function is obtained as

$$\mathcal{F}_{\mathbf{k}; nn'}^<(\omega) = \sum_{lm} G_{\mathbf{k}; n, m}^{\text{R}}(\omega) \Delta_{\mathbf{k}; m-l}^{(s)/(t)} G_{- \mathbf{k}; n', -l}^>(-\omega) + G_{\mathbf{k}; n, m}^<(\omega) \Delta_{\mathbf{k}; m-l}^{(s)/(t)} G_{- \mathbf{k}; n', -l}^{\text{R}}(-\omega). \quad (10)$$

With Eq. (10) we can solve Eq. (9) as a non-hermitian eigenvalue problem. An eigenvalue $\lambda_{\text{sc}} > 1$ signals the onset of superconducting order, and the corresponding eigenvector contains information about the symmetry of the gap function. Detailed derivations are provided in the SM [55].

In the numerical results, the units are set by the hopping $\tilde{t} = 1$, leading to a bandwidth $W \sim 8$. We choose a small bath hybridization $\gamma = 0.1$, set next-to-nearest-neighbor hopping to $\tilde{t}' = -0.2$ and fix the time-averaged filling of the noninteracting (but open and driven) system to $n_f = 0.6$, away from half filling. This is done through the chemical potential of the reservoir μ , where the reservoir temperature is set to $T_{\text{res}} = 0$. We consider a drive frequency regime within the bandwidth $\Omega \leq W$, from low- ($\Omega = 2.0$) to intermediate- ($\Omega = 5.0$) to high-frequency driving ($\Omega = 8.0$) where photoassisted processes become relevant.

Floquet-induced response transition. In Fig. 1, we present the Floquet engineered noninteracting susceptibility and spectral function. For reference, the $\tilde{t}' = 0$ case has been studied

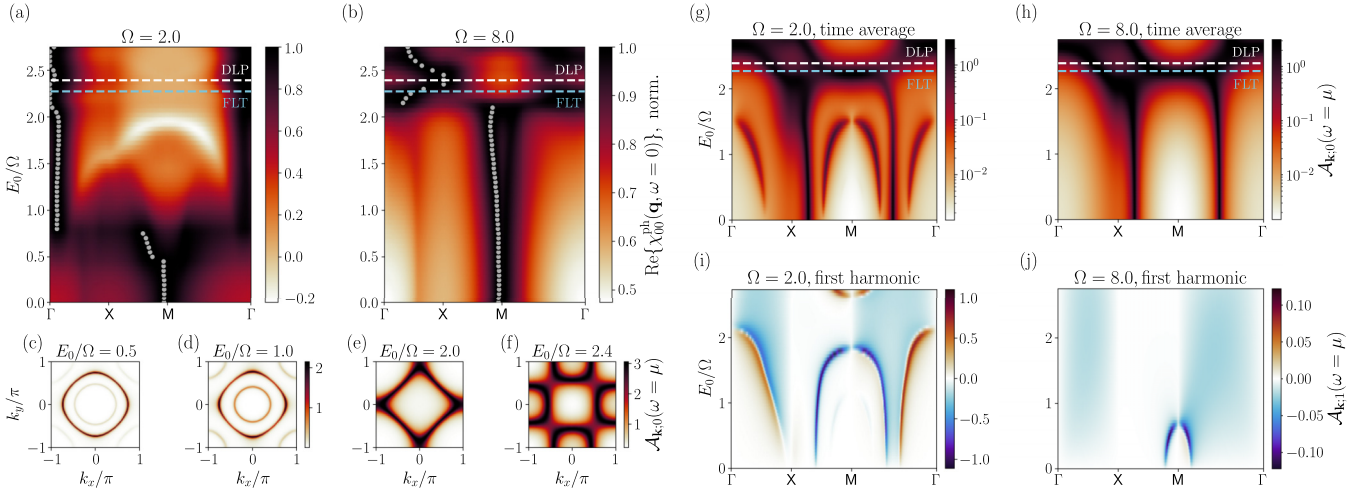


FIG. 1. Time-averaged driven noninteracting susceptibility at $\omega = 0$ and spectral function at $\omega = \mu$. (a), (b) Normalized particle-hole susceptibility in momentum space plotted over drive amplitude. (a) shows low drive frequency, and (b) high drive frequency. Gray dots mark the maximum of the susceptibility for a given amplitude. The blue dashed lines mark the Floquet Lifshitz transition (FLT) and the black dashed lines mark the dynamical localization point (DLP). (c) and (d) show the Fermi surfaces for $\Omega = 2.0$ at the transition point of the susceptibility maximum. (e) and (f) show the same for the high-frequency case ($\Omega = 8.0$, at bandwidth driving). (g) and (h) show the time-averaged spectral function at $\omega = \mu$ in momentum space over drive amplitude for both low and high frequency. Again, the dashed lines mark the Floquet-Lifshitz transition and the dynamical localization point. (i) and (j) show the first harmonic ($n = 1$ Floquet component) of the spectral function.

in Ref. [39]. For $t' \neq 0$, the periodic drive leads to a Floquet Lifshitz transition of the time-averaged Fermi surface [given by the spectral function $\mathcal{A}_{\mathbf{k},0}(\omega = \mu)$] in the vicinity of the dynamical localization point [Figs. 1(e)–1(h)], which also affects the two-particle response function. Figures 1(a) and 1(b) show the static ($\omega = 0$ but periodically driven $\Omega \neq 0$) non-interacting susceptibility evaluated along a path connecting the high-symmetry points of the Brillouin zone as a function of the drive amplitude-to-frequency ratio, E_0/Ω . Figure 1(b) presents the result for the high-frequency driving, $\Omega = 8.0$, where the effects of the Floquet Lifshitz transition and dynamical localization reflect as a transition of the response peak away from the M point towards the X point. Figure 1(a) shows low-frequency driving, $\Omega = 2.0$. Here, the transition occurs at substantially smaller $E_0/\Omega \sim 0.8$. Before this transition, incommensurate peaks emerge at approximately $E_0/\Omega \sim 0.4$, indicated by a small shift of the peak away from the M point. This behavior cannot be explained by either the Floquet Lifshitz transition or dynamical localization, as in the $\Omega = 8.0$ case. Instead, one needs to take into account scattering processes not only within the “main” Fermi surface in the time-averaged component [Figs. 1(c) and 1(d) show no significant change], but also between the Floquet sidebands in the time average and higher harmonics that include photoassisted hoppings and lead to a nonthermal electron occupation; see Figs. 1(g)–1(j) and Ref. [39]. In this way, an additional overlap of spectral weight with almost zero momentum transfer is generated by the drive, the process only being enabled by the scattering with multiple photons.

Floquet engineering magnetic order. The critical interaction strength U_c , above which the system develops magnetic order, is shown in Fig. 2(a) as a function of E_0/Ω . For the high-frequency case, $\Omega = 8.0$, we observe a monotonic decrease of U_c , whose behavior closely follows the results

obtained from a first-order expansion of $\chi_{00}^{\text{ph}}(\mathbf{q} = \mathbf{q}_c, \omega = 0)$, denoted by $\mathcal{O}(E_0/\Omega)$, where $\mathbf{q}_c$ is the momentum giving U_c in the Brillouin zone (see SM for details [55]). This can be interpreted as the peak of the noninteracting susceptibility around the M point increasing in magnitude as $\tilde{t}_{\text{eff}} \rightarrow 0$ without a change in position, as shown in Fig. 1(b) up to $E_0/\Omega \sim 2$. A further increase of E_0/Ω continues to decrease U_c , despite the change in the peak position induced by the Floquet Lifshitz transition. At the dynamical localization point ($\tilde{t}_{\text{eff}} = 0$), the nonzero values of $\tilde{t}'_{\text{eff}}$ and γ prevent U_c from reaching zero. More complex behavior is to be expected in the low-frequency regime, $\Omega = 2.0, 4.0$. We find a suppression of magnetic order that can also be captured by an expansion up to second order, $\mathcal{O}[(E_0/\Omega)^2]$, as denoted by the open dots in Fig. 2(a). This is similar to the results presented in Ref. [34] for the perfect nesting scenario. After the photoassisted response transition, U_c starts to decrease again and the expansion breaks down. In this regime, the strong band deformation dominates over photoassisted processes [see Fig. 1(g) at $E_0/\Omega \gtrsim 1.5$], restoring dynamical localization.

A nonequilibrium phase diagram of the time-averaged steady state is presented in Figs. 2(b) and 2(c), revealing the dominant magnetic phases above the critical interaction U_c [identified in Fig. 2(a)] as well as possible superconducting phases in the vicinity of U_c (see the next section). For both low and intermediate frequencies with $U > U_c$, Néel-type AFM order appears at small E_0/Ω . Upon increasing E_0/Ω , FM order emerges since the peak position of the susceptibility is now located near the Γ point, as shown in Fig. 1(a). At intermediate to high frequencies ($\Omega = 5.0, 8.0$), a stripe AFM ordering develops on the next-to-nearest-neighbor bonds, denoted as AFM2 in Fig. 2(c), at the dynamical localization point. At low-frequency driving, photoassisted processes destroy the AFM2 phase, leaving only the FM phase [Fig. 2(b)].

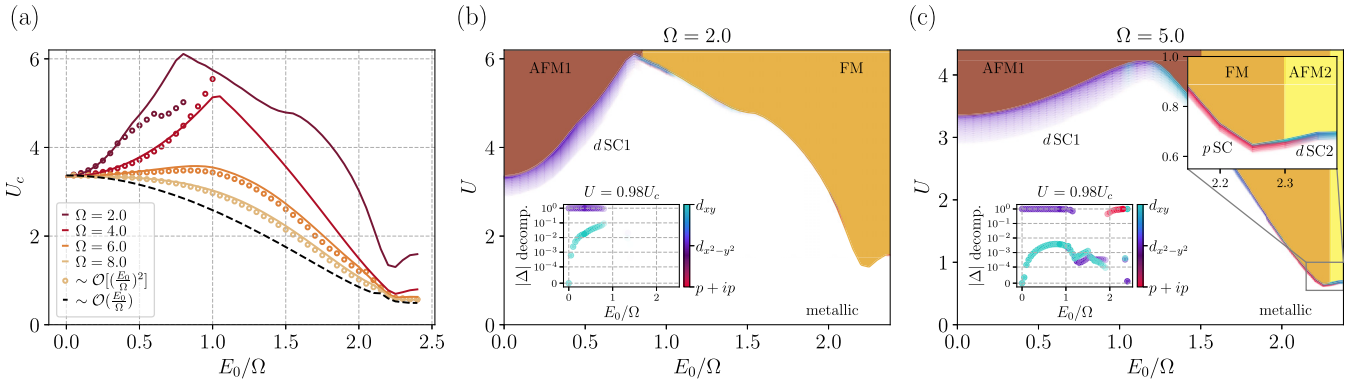


FIG. 2. Interplay of magnetic order and superconductivity in the weak-to-moderate interacting regime. (a) The critical interaction strength for the onset of magnetic order under circularly polarized light as a function of E_0/Ω . Open dots show results obtained by an expansion of $\chi_{00}^{\text{ph}}(\mathbf{q} = \mathbf{q}_C, \omega = 0)$ up to second order $\mathcal{O}[(E_0/\Omega)^2]$ (see SM for details [55]). The Dashed line shows the first-order expansion $\mathcal{O}(E_0/\Omega)$ for $\Omega = 8.0$. (b), (c) Tentative nonequilibrium phase diagram of the time-averaged Floquet steady state at low- and intermediate-frequency driving. AFM and FM order are identified by the position of the peak of the susceptibility in momentum space. AFM1 is conventional Néel-type AFM order on the square lattice, while AFM2 also entails spin antialignment but on the enlarged and tilted lattice given by the next-nearest neighbors (the susceptibility diverges at the X point), leading to stripe ordering. The superconducting regimes $d\text{SC1}$ and $d\text{SC2}$ correspond to the respective AFM transitions and show different mixings of the $d_{x^2-y^2}$ - and d_{xy} -symmetry components. Triplet $p + ip$ -wave superconductivity ($p\text{SC}$) is found close to the FM transition at intermediate drive frequency [(c)]. Insets: Decomposition of the respective symmetry components of the superconducting order.

Superconductivity near the magnetic instability. After establishing the modified magnetic phase diagram of the Floquet-driven Hubbard model, we continue to investigate spin-fluctuation-mediated superconductivity. The possible symmetries of the superconducting order parameter are given by the irreducible representations of the C_{4v} point group of the underlying square lattice. In addition to the s wave, the irreps are given in momentum space by $\Delta_{d_{x^2-y^2}}(\mathbf{k}) = \cos k_x - \cos k_y$, $\Delta_{d_{xy}}(\mathbf{k}) = \sin k_x \sin k_y$, and $\Delta_{p_{x(y)}}(\mathbf{k}) = \sin k_{x(y)}$. By projecting the superconducting gap onto the irreps we can determine the gap symmetry [41]. In Figs. 2(b) and 2(c), one can find different superconducting regimes in the vicinity of magnetic order for the low- and intermediate-frequency cases. In the absence of driving ($E_0/\Omega = 0$), but with coupling to the reservoir, we find spin-fluctuation-induced $d_{x^2-y^2}$ pairing, as expected from equilibrium [62–64]. Upon turning on the drive, the $d_{x^2-y^2}$ ordering persists in the vicinity of Néel-type AFM ordering (AFM1) for both low ($\Omega = 2.0$) and intermediate ($\Omega = 5.0$) frequencies up to $E_0/\Omega \lesssim 1$. Although the dominant pairing symmetry remains $d_{x^2-y^2}$, driving induces a mixing with d_{xy} symmetry components (see the inset). To investigate the mixing of the subcomponents and their topological properties, we calculate the Berry curvature of the associated Bogoliubov–de Gennes Hamiltonian using Fukui’s method [65]. We find that a topological $d_{x^2-y^2} + id_{xy}$ order with Chern number $C = 2$ is induced by driving in the vicinity of AFM1 for both $\Omega = 2.0, 5.0$. In the low-frequency case [Fig. 2(b)], chiral d -wave superconductivity is suppressed as soon as the magnetic instability switches to FM order. Although triplet superconductivity could be expected to emerge in the vicinity of the FM instability, we do not find this order for the given parameters. However, in the intermediate-frequency regime [Fig. 2(c)], triplet $p + ip$ -wave superconductivity can be stabilized around $E_0/\Omega \sim 2$ close to

the Floquet Lifshitz transition [see also the inset in Fig. 2(c)]. With a further increase of E_0/Ω , the effect of the dynamical localization point becomes evident, where $d_{x^2-y^2} + id_{xy}$ order is stabilized again, but now with the d_{xy} component being dominant ($d\text{SC2}$). This phase appears in the vicinity of the magnetic instability towards the stripe AFM order (AFM2), rather than AFM1. All orders are found to be topologically nontrivial (see SM [55]).

Discussion. We leveraged methodological advances to theoretically describe magnetic order and spin-fluctuation-mediated superconductivity in the circularly driven square-lattice Hubbard model. Introducing the Floquet random-phase approximation approach we could study the effective pairing vertex and the resulting superconducting gap under subband-width driving. At low drive frequency, photoassisted hopping and the resulting nonthermal electron occupation generally suppress magnetic order, whereas at strong drive amplitude, band deformation can still lead to a dynamical localization effect that enhances magnetic order. However, the Floquet engineered fluctuations change the ordering tendency from AFM to FM order. We then investigated the fate of superconducting order in the vicinity of the magnetic instability, finding that superconductivity can generally survive and is engineered into topological $d + id$ order that is stable over a wide range of drive parameters. While low-frequency driving with high amplitude towards the FM regime destroys any superconducting order, at intermediate-frequency driving, p -wave superconductivity emerges together with a second type of topological $d + id$ -wave ordering in the vicinity of the Floquet Lifshitz transition.

We have shown that analytical expansions cannot be reliably extrapolated from their respective limits, but instead the complex interplay of multiphoton processes and deformed bands needs to be fully accounted for in the

weak-to-moderate interacting regime. Further improvements to the method would involve its extension to truncated-unity functional renormalization group [66,67] and to treat inelastic scattering processes by including frequency and momentum degrees of freedom simultaneously. Both avenues require further advances in quantum field theory simulations, a promising idea being, for example, tensor cross interpolation [68–70].

Our theoretical simulations highlight parameter regimes that are accessible in pulsed experiments [8] in which Floquet engineering allows to control the subtle interplay of magnetic and superconducting order. This offers different pathways for understanding the nature of unconventional superconductivity in cuprates [15], rhombohedral multilayer graphenes [71–73], or moiré materials [74–77].

Acknowledgments. We acknowledge fruitful discussions with Qiyu Liu. J.-N.H. acknowledges funding by the Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) - 508440990 - 531215165 (Research Unit “OPTIMAL”). T.O. acknowledges support from the JSPS Overseas Research Fellowships. Simulations were performed with computing resources granted by RWTH Aachen University under project rwth1749 as well as the HPC system Raven and Viper at the Max Planck Computing and Data Facility.

Data availability. The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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