

Resonant false vacuum decay in two dimensions on a 4000-qubit quantum annealer

Jaka Vodeb

`jaka.vodeb@ijs.si`

Jozef Stefan Institute <https://orcid.org/0000-0002-6193-3475>

Gregor Humar

Department of Complex Matter, Jozef Stefan Institute, Jamova 39, 1000 Ljubljana, Slovenia;
Department of Physics, Faculty for Mathematics and Physics, Jadranska 19, University of Ljubljana, Slovenia <https://orcid.org/0009-0003-8551-3920>

Jean-Yves Desaulles

Institute of Science and Technology Austria (ISTA) <https://orcid.org/0000-0002-3749-6375>

Luka Pavešić

University of Padova

Marko Ljubotina

Technical University of Munich, TUM <https://orcid.org/0000-0003-0038-7068>

Zlatko Papic

University of Leeds <https://orcid.org/0000-0002-8451-2235>

Kristel Michielsen

Forschungszentrum Jülich GmbH <https://orcid.org/0000-0003-1444-4262>

Article

Keywords:

Posted Date: August 26th, 2026

DOI: <https://doi.org/10.21203/rs.3.rs-10244511/v1>

License:   This work is licensed under a Creative Commons Attribution 4.0 International License.
[Read Full License](#)

Additional Declarations: There is **NO** Competing Interest.

Resonant false vacuum decay in two dimensions on a 4000-qubit quantum annealer

Gregor Humar,^{1,2} Jean-Yves Desaulles,³ Luka Pavešić,^{4,5} Marko Ljubotina,^{6,7} Zlatko Papić,⁸ Kristel Michielsen,^{9,10} and Jaka Vodeb^{1,11,12,*}

¹*Department of Complex Matter, Jožef Stefan Institute, Jamova 39, 1000 Ljubljana, Slovenia*

²*Department of Physics, Faculty of Mathematics and Physics, University of Ljubljana, Jadranska 19, 1000 Ljubljana, Slovenia*

³*Institute of Science and Technology Austria (ISTA), Am Campus 1, 3400 Klosterneuburg, Austria*

⁴*Department of Physics and Astronomy “G. Galilei”, University of Padova, via Marzolo 8, I-35131 Padova, Italy*

⁵*National Institute for Nuclear Physics (INFN), Section of Padova, I-35131 Padova, Italy*

⁶*Technical University of Munich, TUM School of Natural Sciences, Physics Department, James-Frank-Str. 1, 85748 Garching, Germany*

⁷*Munich Center for Quantum Science and Technology (MCQST), Schellingstr. 4, München 80799, Germany*

⁸*School of Physics and Astronomy, University of Leeds, Leeds LS2 9JT, UK*

⁹*Jülich Supercomputing Centre, Forschungszentrum Jülich, D-52425 Jülich, Germany*

¹⁰*Department of Computer Science, University of Cologne, 50931 Cologne, Germany*

¹¹*CENN Nanocenter, Jamova 39, 1000 Ljubljana, Slovenia*

¹²*Jülich Supercomputing Centre, Institute for Advanced Simulation, Forschungszentrum Jülich, 52425 Jülich, Germany*

(Dated: July 3, 2026)

From cosmology to quantum matter, metastable states often decay through the nucleation and growth of competing domains, with false vacuum decay providing the paradigmatic example of this process. Here we demonstrate a distinct regime in which domain growth outpaces nucleation by orders of magnitude and is controlled by local resonance conditions. Using a programmable quantum annealer with more than 4000 qubits, we realize a two-dimensional quantum Ising model whose metastable spin-polarized state encodes a false vacuum. At a specific value of the longitudinal field, single-spin flips at the boundary of a seeded bubble become resonant, enabling kinetically constrained expansion. Combining experiment with tensor-network simulations and stochastic circuit modeling, we observe nearly ballistic growth of true-vacuum domains with sub-ballistic interface broadening, consistent with Kardar–Parisi–Zhang universality. Our results establish a growth-dominated regime of false vacuum decay and show how large-scale quantum simulation can access nonequilibrium metastable dynamics relevant to quantum field theory, cosmology, and strongly correlated matter.

INTRODUCTION

Metastability is ubiquitous in many-body physics: a state can remain long-lived even when a lower-energy configuration is available, until fluctuations eventually trigger its decay. In quantum field theory and cosmology, the relaxation of a metastable vacuum through the nucleation of localized bubbles of the true vacuum is known as false-vacuum decay (FVD) [1–4]. More broadly, the same phenomenon underlies first-order quantum phase transitions [5] and the decay of metastable states in settings ranging from mesoscopic systems and quantum materials to reaction-rate theory, open quantum systems, and many-body quantum dynamics [6–11].

The conventional description of FVD is critical-bubble theory: quantum or thermal fluctuations nucleate small domains of the stable phase, and the competition between bulk energy gain and interfacial tension determines whether such bubbles shrink or expand [1, 2, 4]. Once a bubble exceeds a critical size, its subsequent evolution is usually regarded as comparatively simple, with growth driven by the free-energy difference between the two phases. In this way, standard treatments separate the problem into *nucleation* and *growth*, with most of the conceptual emphasis placed on the former [12–24]. This leaves open the possibility of a qualitatively different regime in which post-nucleation growth itself becomes the nontrivial part of the dynamics, controlled by local resonant processes at the bubble boundary rather than by the smooth pressure-driven motion of an interface.

Motivated by the field-theoretic picture, lattice realizations of FVD in quantum spin systems were first explored theoretically [25–28] and have recently become experimentally accessible. Analogs of FVD have now been realized in superfluids [29], neutral-atom systems [30, 31], cold atoms [32, 33], trapped ions [34], quantum annealers [35], and

* jaka.vodeb@ijs.si

graphene multilayers [36]. In one dimension, these studies probed decay-rate scaling, confinement effects, and the importance of properly preparing metastable states [26, 27, 31], revealing intrinsic lattice effects, such as quantized and resonant nucleation channels [31, 35]. By contrast, FVD in two dimensions remains far less understood, with recent numerical studies starting to access such regimes [37] and Rydberg atom experiments observing resonant deconfined cluster-growth dynamics governed by kinetic constraints [38], which we also address in this work. In particular, the real-time growth of resonant false-vacuum bubbles in two dimensions has largely remained unexplored.

Two dimensions are qualitatively different from one: bubble growth is no longer constrained to a line, but depends on the geometry of the interface and on the competition between area and perimeter. This makes two dimensions a natural setting to explore whether post-nucleation growth can itself become the essential part of FVD. Here, we address this question using a large-scale programmable quantum annealer that realizes the two-dimensional mixed-field Ising model in a metastable regime. We identify a special resonance at which single-spin flips at the boundary of an existing true-vacuum seed become energy conserving, causing bubble growth to outpace spontaneous nucleation by several orders of magnitude. Combining quantum simulation with tensor-network calculations and large-scale stochastic modeling, we show that this produces a distinct growth-dominated regime of FVD, characterized by ballistic expansion together with nontrivial interface broadening following the Kardar-Parisi-Zhang (KPZ) universality class. Our results uncover a new route to FVD in which local resonant processes govern the expansion of true-vacuum domains, and establish programmable quantum hardware as a controlled platform for studying FVD beyond the standard semiclassical paradigm.

QUANTUM SIMULATION OF RESONANT FALSE VACUUM DECAY

We study the mixed-field quantum Ising model

$$\hat{H} = -J \sum_{\langle i,j \rangle} \hat{\sigma}_i^z \hat{\sigma}_j^z - h_x \sum_i \hat{\sigma}_i^x - h_z \sum_i \hat{\sigma}_i^z, \quad (1)$$

on a square lattice with periodic boundary conditions, in the regime $J \gg h_x$. For $h_z = 0$, the model has two degenerate ferromagnetic ground states, $|\downarrow\downarrow\ldots\downarrow\rangle$ and $|\uparrow\uparrow\ldots\uparrow\rangle$. A finite longitudinal field lifts this degeneracy, so that one state becomes the true vacuum and the other a metastable false vacuum; here we take the all-down state as the true vacuum and the all-up state as the false vacuum. Figures 1a,b show the implementation of this model on the D-Wave Advantage2 quantum annealer. The logical lattice has the topology of a torus and is embedded into the Advantage2_system1.6 and Advantage2_system1.13 processors, using 4058 and 4044 out of 4593 and 4579 physical qubits, respectively. The reverse-annealing protocol provides direct control over the time-dependent transverse field $h_x(t)$, while the longitudinal field $h_z(t)$ can be programmed independently, as described below. The exact details of the protocols are given in the Supplementary Information (SI) [39].

Domains of opposite magnetization of different shapes and sizes in the background of the metastable polarized state are true-vacuum bubbles. The energy of a domain with area A and perimeter ∂A is determined by the negative bulk contribution $2h_z A$ and the positive interface term $2J \partial A$. For specific values of h_z and J , certain domains become resonant with the initial metastable state. Resonances occur when

$$2|h_z|A = 2J \partial A. \quad (2)$$

As shown in Fig. 1c, this gives one-spin, two-spin, three-spin, and symmetric four-spin resonances at $h_z = 4J, 3J, 8J/3$, and $2J$, respectively. Their amplitudes scale parametrically as h_x^A , so the one-spin resonance is the strongest, followed by the two- and three-spin resonances. This hierarchy is directly visible by holding h_z constant at various magnitudes while ramping $h_x(t)$ from 0 to a small value and back in Fig. 1d, where the strongest features correspond to the lowest-order processes, and the snapshots on the right confirm the associated bubble geometries.

The resonance at $h_z = 2J$ is special because it not only allows four-spin nucleation, but also first-order growth of an already existing bubble [38]. For an up-spin adjacent to a down-spin seed, the local Ising contribution from its neighborhood is $2J$, so the effective longitudinal term becomes $-\hat{\sigma}_i^z(2J - h_z)$. Geometrically, the move increases the bubble area by one site while creating two additional domain-wall segments, increasing the energy by $4J$. This increase is exactly offset by the longitudinal field at $h_z = 2J$, making the growth process resonant. Once a seed is present, the bubble can therefore spread through single-spin flips, a first-order process in h_x . In comparison, spontaneous nucleation of the four-spin bubble is fourth order, so growth dominates over nucleation by approximately three orders of magnitude in the experimentally relevant regime.

To probe the growth regime directly, we prepare a single flipped spin at the center of the system and ramp the global h_z through the $h_z = 2J$ resonance while h_x remains finite, to isolate the fast interface dynamics. The time spent in the resonant regime is estimated to be of order tens of nanoseconds, consistent with the coherent operating

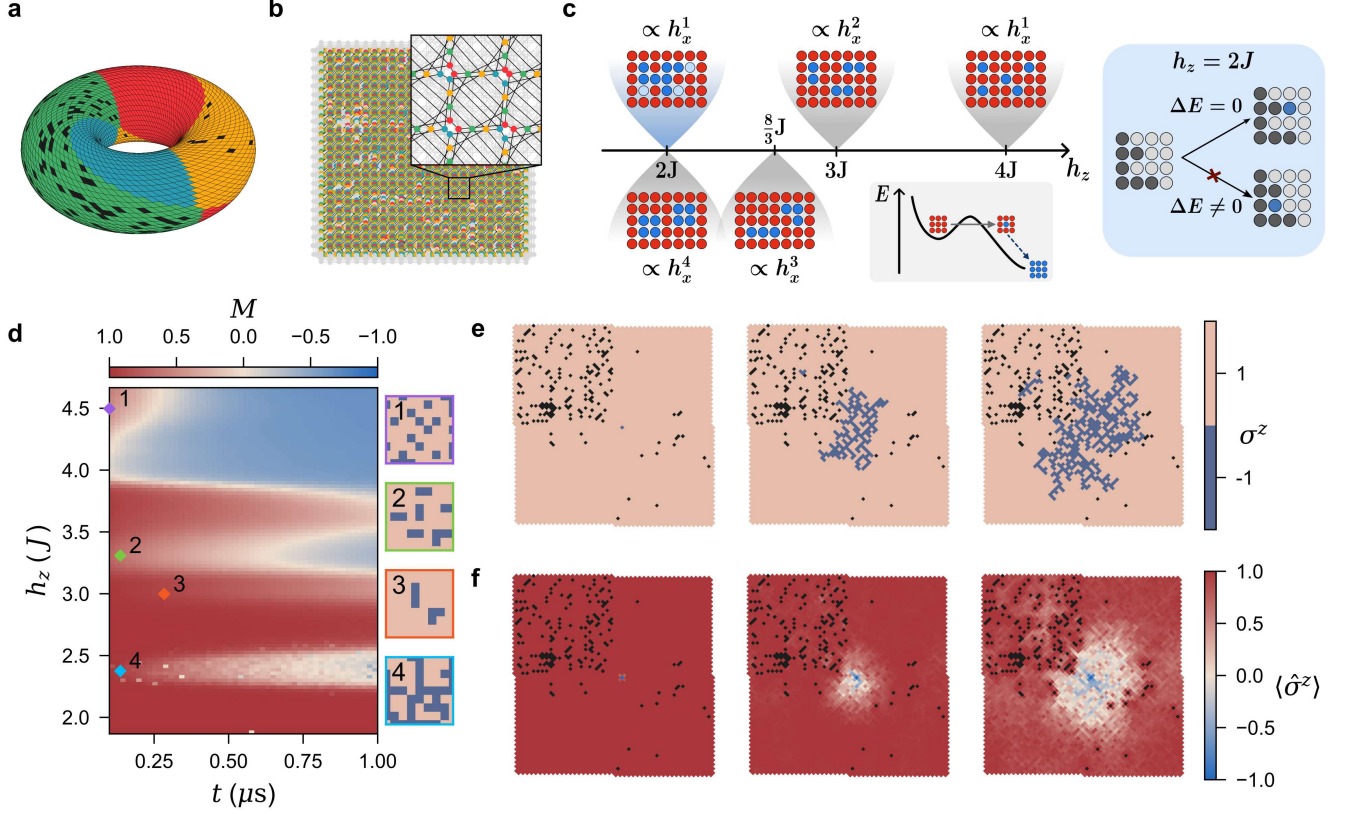


FIG. 1. Two-dimensional resonant false-vacuum decay on a quantum annealer. **a**, Logical square lattice with periodic boundary conditions in both directions, realizing a torus topology. The colors indicate subsets of logical sites according to the physical-qubit groups used in the embedding shown in **b**. Black squares denote missing physical qubits on the processor, which are bypassed in the logical construction. **b**, Embedding of the logical lattice on the quantum processing unit. The colored groups match those in **a** and are connected so as to implement the periodic two-dimensional geometry despite hardware defects. **c**, Lowest resonances of the mixed-field Ising model as the longitudinal field is varied. The resonances correspond to the condition in Eq. 2 where the domain-wall cost matches the energy gain from flipped spins. Shown are the one-, two-, three-, and four-spin resonances. The $h_z = 2J$ resonance is special because, in addition to four-spin nucleation, it also enables first-order resonant growth of an already existing bubble. **d**, Experimental sweep of the longitudinal field on the annealer, showing the magnetization M as a function of h_z and evolution time t for the constant- h_z reverse-annealing protocol. The measured resonance structure agrees with the sequence shown in **c**; representative spin configurations at the numbered points are displayed on the right. The strongest resonances are the lowest-order ones, while the higher-order two- and three-spin resonances appear progressively weaker, consistent with the hierarchy of tunneling amplitudes. The resonant values of h_z are slightly shifted with respect to expected values ($4J$, $3J$, $8J/3$, $2J$) due to slight miscalibration between the user specified values of h_z and J and actual device values, however, we characterized the resonances using the dominant bubble type we observed. **e**, Three single-shot experimental configurations obtained from a deliberately prepared initial state containing a single flipped spin at the center of the system. Here, a time-dependent linear ramp of $h_z(t)$ is used to pass through the $h_z = 2J$ resonance while h_x remains finite, thereby isolating the fast growth dynamics associated with the resonant interface process. A ramified fractal pattern spreads outward from the seed in a manner consistent with the local resonance rule. At later times, weak thermalization effects become visible through violations of the strict resonance constraint. **f**, Time-averaged spin configurations obtained from 100 snapshots of the type shown in **e**, revealing the average spatial structure of the resonantly spreading bubble.

regime established in previous annealer-based quantum simulations [40–42]. Further details are given in the SI [39]. Representative single-shot configurations are shown in Fig. 1e. Starting from the seeded defect, the true-vacuum region expands in a ramified, fractal-like pattern consistent with the local resonance rule: only interface spins in the proper environment can flip resonantly. At later times, weak violations of this ideal pattern appear, signaling the onset of off-resonant processes, but the dynamics remain dominated by the resonant growth channel. Averaging over many realizations, Fig. 1f reveals a growing, approximately isotropic domain, centered on the seed. Thus, the $h_z = 2J$ resonance gives rise not merely to isolated resonant moves, but to a robust growth regime with an emergent propagating interface.

TRUE VACUUM BUBBLE GROWTH

We further characterize the true vacuum bubble growth obtained using the ramp protocol with an artificial initial seed at multiple locations, used for averaging. As shown in Fig. 2a, the local magnetization profile, $\langle \hat{\sigma}_j^z \rangle$, measured along a horizontal cut through the center, develops a clear light-cone structure. The initially localized seed spreads outwards approximately symmetrically, forming an expanding true-vacuum domain inside the false-vacuum background. Figure 2b shows the corresponding nearest-neighbor kink density $\bar{K}_{j,j+1} = (1 - \hat{\sigma}_j^z \hat{\sigma}_{j+1}^z)/2$. Its sharp expanding front matches the light-cone in the magnetization which demonstrates that the growth is controlled by the motion of a well-defined interface. For a standard expanding bubble, one would expect the kink density in the bulk to decrease to zero. However, we find a large kink density in the bulk, suggesting a rich internal structure of the growing false vacuum domain, consistent with the ramified pattern.

To characterize the propagation more quantitatively, we study the equal-time correlation function $\langle \hat{\sigma}_{-j}^z \hat{\sigma}_j^z \rangle$ and its connected counterpart $\langle \hat{\sigma}_{-j}^z \hat{\sigma}_j^z \rangle_c = \langle \hat{\sigma}_{-j}^z \hat{\sigma}_j^z \rangle - \langle \hat{\sigma}_{-j}^z \rangle \langle \hat{\sigma}_j^z \rangle$, where $\hat{\sigma}_j^z$ denotes the single-site magnetization at horizontal distance j from the central initialized bubble site. The resulting space-time maps are shown in Fig. 2c,d. In both cases, the correlation front propagates linearly to a good approximation, allowing us to define an edge position as a function of time. Fitting the extracted front position to a power law $c\tau^p$ gives exponents $p \approx 1.15 \pm 0.06$ for the full correlator and $p \approx 1.00 \pm 0.10$ for the connected correlator. Repeating the same analysis for different values of h_x , Fig. 2e shows that the extracted exponent remains close to $p = 1$ across the explored range. The larger deviations at small h_x arise from slower dynamics, where the corresponding changes in the magnetization profile are harder to resolve and increase the fit uncertainty. By contrast, at large h_x , the fitted exponent is systematically shifted above $p = 1$ because the bubble reaches the system boundary and self-interferes through the periodic boundary conditions. These data indicate that the resonant bubble grows nearly ballistically in the experimentally accessible time window.

This conclusion is further illustrated by the magnetization profiles in Fig. 2f. As time increases, the front moves to larger distances while retaining an approximately self-similar shape. We find a collapse by rescaling the coordinate $j\tau^{-1}$ in Fig. 2g, where τ is the duration of the linear $h_z(t)$ ramp and is proportional to the time spent in the resonant regime, confirming that the dominant scaling is ballistic. At the same time, the residual deviations from perfect collapse near the interface indicate that the front is not rigid. We address this observation in the next section.

The results presented here are expected to generalize to larger scales and higher dimensions in a weak transverse field regime $h_x \ll J$. In $d = 2$ dimensions, similar resonances occur at $|h_z| = 2J/l$, where filaments with thickness l spread ballistically from an initial seed of size $4l^2$ with velocity h_x^l . The rate of initial seed nucleation scales with its volume ($h_x^{4l^2}$), while the rate of filament spreading scales with the square root of the volume (h_x^l), thereby governing the late-time dynamics. A similar resonant phenomenon is also expected in $d > 2$, where the analogous symmetric initial seed would be a (hyper-)cube with side $\frac{d}{d-1}l$ and the nucleation rate scaling as $\sim h_x^{\left(\frac{d}{d-1}l\right)^d}$. The corresponding resonances appear at $|h_z| = 2J(d-1)/l$ where filaments with thickness l would grow ballistically with the rate that scales as $\sim h_x^{l^{d-1}}$, still dominating the dynamics in any d . The ratio of nucleation and filament growth rate exponents is $l \left(\frac{d}{d-1}\right)^d$, implying that filament growth becomes even faster compared to nucleation with increasing dimensionality.

CLASSICAL MODELING OF BUBBLE GROWTH

We support the quantum simulation results with two complementary classical approaches. First, we simulate the full quantum model of Eq. (1) with tree tensor networks (TTNs), a generalization of matrix product states that better represent long-range correlations in higher-dimensional systems [43]. TTNs allow us to simulate the full quantum dynamics with controlled truncation errors, albeit with limitations on the accessible system sizes and total evolution times [44, 45]. To access the late-time properties of bubble growth beyond the reach of TTNs, we additionally develop an effective stochastic circuit (SC) model of bubble dynamics. In the SC approach, we initialize a classical configuration with a single lattice site flipped into the true-vacuum state and evolve it in discrete time steps denoted by t_s . At each step, any site j is acted on with a constrained spin-flip operator, preserving the resonance condition, with probability p and left unchanged with probability $1 - p$. The crucial point is that each update maps a single computational-basis configuration to another computational-basis configuration, and can thus be simulated classically. Further details about the two types of classical simulations can be found in Methods.

To isolate the intrinsic bubble-growth dynamics from the extraneous ramp protocol effects present in quantum simulations, we consider instantaneous quenches from an initial product state in which a single spin is flipped into the true-vacuum state. We then evolve under the full Hamiltonian at finite h_x and fixed $h_z = 2J$. The evolution of the magnetization, computed by TTNs, is shown in Fig. 3a, while a representative trajectory of the SC simulation is shown in the top three panels of Fig. 3b. Each trajectory exhibits the same tree-like pattern reminiscent of the

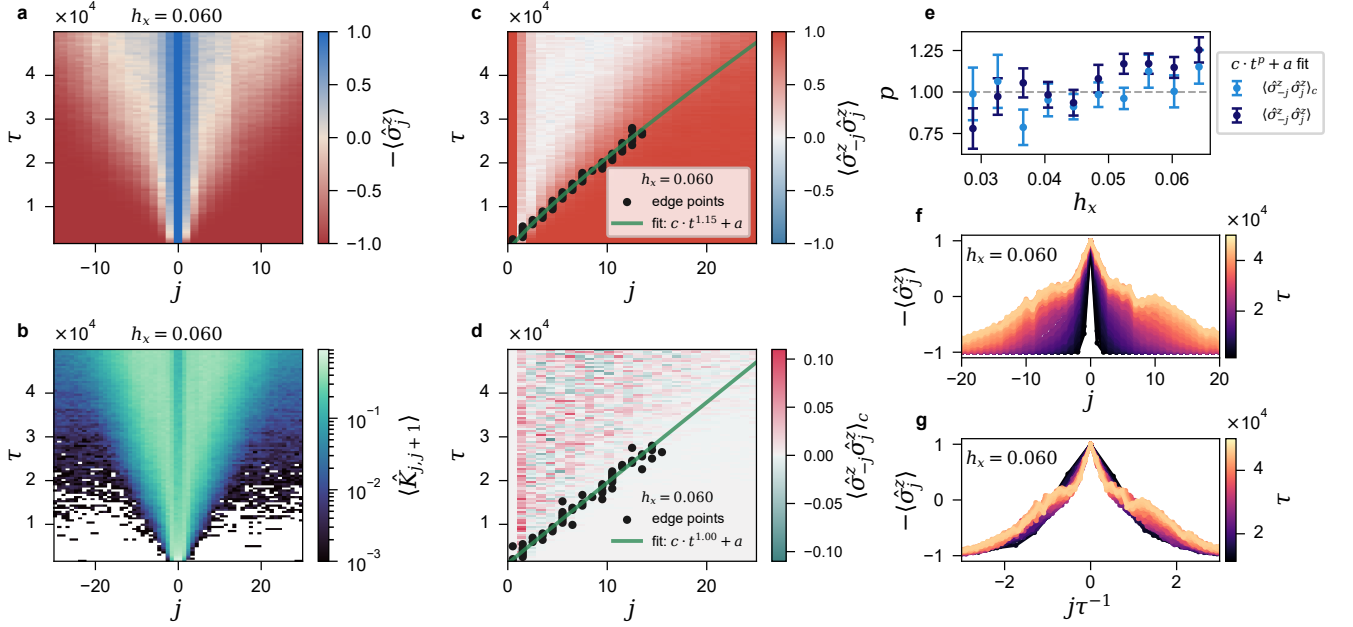


FIG. 2. **Experimental light-cone dynamics of resonant bubble spreading on the quantum annealer.** **a**, Time evolution of the measured local magnetization profile $-\langle\hat{\sigma}_j^z\rangle$ along a horizontal cut through the center of the two-dimensional system, shown as a function of distance j from the initially prepared central seed and evolution time τ , for $h_x = 0.06$. The bright central region corresponds to the expanding true-vacuum bubble, whose front forms a clear light-cone structure. **b**, Corresponding nearest-neighbor kink density $\langle\hat{K}_{j,j+1}\rangle$, showing the propagation of the domain-wall fronts that bound the bubble. The two fronts move approximately symmetrically away from the seeded center and directly visualize the spreading interface. **c**, Equal-time correlation function $\langle\hat{\sigma}_{-j}^z\hat{\sigma}_j^z\rangle$, where $\hat{\sigma}_j^z$ denotes the single-site magnetization at horizontal distance j from the central initialized bubble site. The black points mark the extracted edge positions of the correlation front, and the green line shows a power-law fit $c \cdot \tau^p$ with exponent $p \approx 1.15$, indicating nearly ballistic spreading with a weak apparent superlinear correction over the experimentally accessible times. **d**, Connected equal-time correlator $\langle\hat{\sigma}_{-j}^z\hat{\sigma}_j^z\rangle_c$. The black points again mark the extracted edge positions, while the green line is a power-law fit $c \cdot \tau^p$ with exponent $p \approx 1.00$. The close agreement with **c** shows that the front dynamics is robust with respect to disconnected background contributions. **e**, Growth exponent p extracted from power-law fits to the edge positions of $\langle\hat{\sigma}_{-j}^z\hat{\sigma}_j^z\rangle$ and $\langle\hat{\sigma}_{-j}^z\hat{\sigma}_j^z\rangle_c$ as a function of h_x . The dashed horizontal line marks purely ballistic propagation, $p = 1$. Both correlators yield exponents close to unity across the explored range of transverse fields, consistent with nearly ballistic spreading. **f**, Raw magnetization profiles $-\langle\hat{\sigma}_j^z\rangle$ at different times τ , showing the outward propagation of the bubble front. The approximately triangular profile in the laboratory frame reflects a front moving with nearly constant velocity while broadening in time. **g**, The same data as in **f**, replotted against the scaled coordinate $j\tau^{-1}$. The approximate collapse confirms ballistic scaling of the front position, while the residual deviations near the interface encode the slower broadening of the wavefront.

experimental data in Fig. 1.

The bottom panel of Fig. 3a reveals the breakdown of TTN calculations: the bubble seemingly grows faster in certain preferred directions. Such anomalous interface effects are a consequence of mapping the 2D system into a tensor network that does not natively represent the lattice geometry [46]. This is a sign that the simulation becomes unreliable for $th_x \gtrsim 3$ for this system of 16×16 spins. On the other hand, the SC dynamics in Fig. 3b provides a window to the coarse-grained description of the long-time growth regime, where the dynamics is dominated by the accumulation of many local resonant events and the relevant observables are statistical properties of the growing bubble. This allows us to access the statistical growth dynamics of bubbles in systems containing hundreds of thousands of spins and over thousands of time steps. The magnetization, averaged over many SC realizations, is shown in the bottom panel of Fig. 3b. Interestingly, it demonstrates a remarkable spherical invariance despite the underlying D_4 symmetry of the square lattice.

The evolution of the bubble size r , computed via TTNs from the second moment of $\langle\hat{\sigma}^z\rangle$, is shown in Fig. 3c for a few values of h_x . This data suggests a ballistic growth with a characteristic timescale of h_x , indicating that the growth process is driven by the first-order resonance, rather than by quantum coarsening or other diffusive processes. To extract the long-time features of the growth, we focus on the SC simulation of the same quantity r (see Methods). Figure 3d plots the log derivative of r over time, displaying close agreement and convergence towards the ballistic

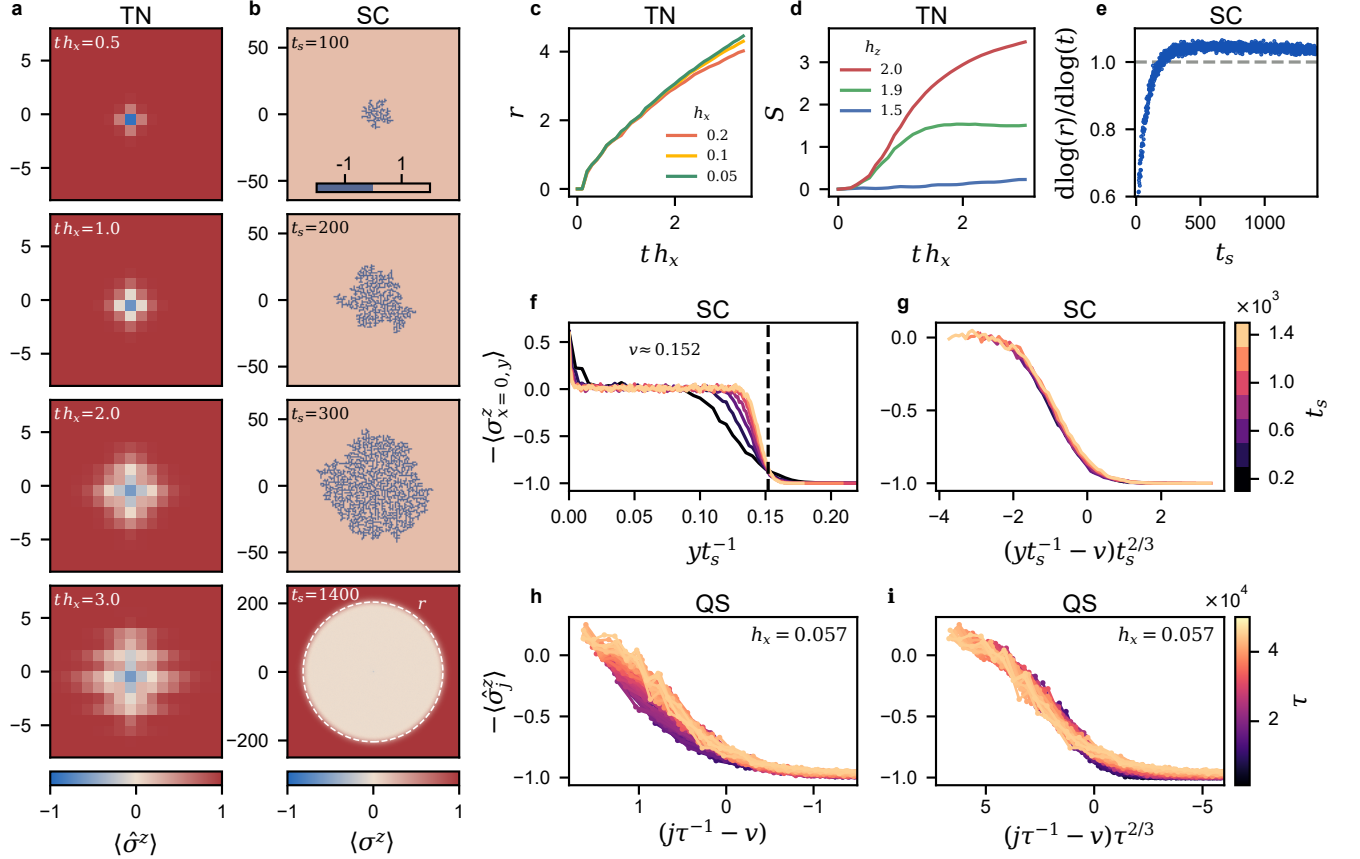


FIG. 3. **Ballistic growth and universal interface broadening of resonantly spreading bubbles.** **a**, TTN simulations of the magnetization $\langle \hat{\sigma}^z \rangle$, showing an expanding true-vacuum domain with a broadening front. **b**, SC simulations of the effective resonant model showing representative snapshots of the magnetization σ^z at increasing times, with the averaged local magnetization $\langle \sigma_{x,y}^z \rangle$ at long time shown at the bottom. The dashed circle has radius r equal to $\sqrt{2}$ times the second moment of the bubble profile, demonstrating that the expanding domain remains approximately isotropic. **c**, Bubble radius extracted from the TTN simulations, plotted as r versus the scaled time th_x . The data collapse for several values of h_x confirms that the growth timescale is given by th_x , as expected for the first-order resonant growth process generated by the effective Hamiltonian. **d**, Logarithmic derivative $d \log(r)/d \log(t)$ extracted from the SC simulations. The effective exponent remains close to unity over the accessible time window and decreases only weakly at late times, consistent with asymptotically ballistic radial growth. **e**, Bipartite entanglement entropy S from the TTN simulations as a function of th_x . The entropy growth remains moderate over the accessible time window, explaining why tensor-network simulations can faithfully capture the early-time quantum dynamics while becoming limited at later times and larger systems. **f**, Collapse of the SC magnetization profile along the line $x = 0$, plotted as $-\langle \sigma_{x=0,y}^z \rangle$ versus y/t . The good overlap of the profiles at different times demonstrates a well-defined ballistic front propagating with velocity $v \approx 0.152$, while deviations near the front indicate additional broadening beyond a rigidly translating profile. **g**, Scaling collapse of the SC wavefront after subtracting the ballistic drift, plotted as a function of $(y/t - v)t^{2/3}$. The collapse is consistent with front broadening governed by the Kardar-Parisi-Zhang (KPZ) scaling exponent $2/3$. **h**, Experimental quantum-simulator (QS) profiles along a one-dimensional cut through the two-dimensional bubble at several evolution times, plotted in the co-moving coordinate $(j/\tau - v)$. The measured magnetization shows the same ballistic front propagation as in the SC and TN calculations. **i**, Collapse of the experimental QS wavefront using the KPZ-scaled coordinate $(j/\tau - v)\tau^{2/3}$. The agreement with the stochastic-circuit scaling supports a universal coarse-grained description of the resonant bubble interface, with nearly ballistic radial growth and sub-ballistic KPZ-type front broadening.

growth exponent of 1.

Despite the fact that the main features of the bubble growth can be reproduced by SC simulations, the underlying dynamics generate a considerable amount of entanglement and should thus be understood as a genuinely many-body quantum process. Figure 3e shows the evolution of the half-cut entanglement entropy S for quenches of h_z to the values indicated in the legend, computed via TTNs. As expected, the state remains approximately frozen when there are no resonant processes available ($h_z = 1.5J$) [47], while the resonant process at $h_z = 2J$ triggers very fast entanglement growth. Interestingly, for a quench close to the resonant point ($h_z = 1.9J$), S initially grows at the

same rate as in the resonant case, but quickly tapers off and saturates at a smaller value. This might be due to the confinement of excitations, expected at small h_x far from the critical point [37, 48].

It is inherently difficult to directly compare TTN results with our quantum simulations because the precise $h_z(t)$ and $h_x(t)$ ramps are unknown, even though they have a strong influence on the exact shape of $M(t)$. However, we consistently find a collapse of $M(t)$ with a rescaled time axis $h_x^2 t$ in quantum simulation, contrary to the expected $h_x t$ scaling. We therefore performed analogous $h_z(t)$ ramps through the resonance using TTNs and found that the $h_x^2 t$ scaling law emerges there as well, see SI [39]. This is explained by the quantum nature of the resonance, where perturbation theory predicts a finite resonance width $\delta h_z \propto h_x$. Then the time spent in the resonance during the ramp is $\delta t \propto h_x$, and the scaling law of the dynamics needs to be appropriately modified as $M(t) \sim h_x \delta t \propto h_x^2 t$, suggesting agreement between TTN and quantum simulation.

An alternative way to test for the ballistic bubble growth is to collapse the data along a single row or column, as is done in Fig 3f. Assuming a constant propagation speed leads to a clean collapse everywhere except near the center of the system and around the wavefront. This indicates that the latter broadens more slowly than ballistically. Performing a second collapse of the wavefront data, Fig 3g, we find that the best results are obtained for $t_s^{2/3}$. This strongly suggests that our stochastic circuit belongs to the KPZ universality class, which describes transport properties of a variety of quantum systems [49–51]. This class was first developed to describe the growth of interfaces in classical systems, which is indeed similar to the process at play here [52]. We perform the same collapses on the quantum simulation data, shown in Fig. 3h,i, and find a good agreement both for the ballistic growth of the bubble size and the KPZ coarsening of the interface, see the SI [39] for a detailed analysis of the optimal exponent across the entire quantum simulation dataset.

DISCUSSION AND OUTLOOK

We have studied a growth-dominated regime of FVD in two dimensions on a superconducting quantum annealer. At the resonance $|h_z| = 2J$, single-spin flips at the boundary of an existing true-vacuum seed become energy conserving, so growth already proceeds at first order in h_x , while spontaneous nucleation of new bubbles remains parametrically slower. This produces an approximately three orders of magnitude separation between growth and nucleation. In our quantum simulations, the resulting dynamics appear as a clear light-cone structure and nearly ballistic spreading of a seeded bubble, while classical modeling captures the accompanying sub-ballistic interface broadening consistent with KPZ scaling. The essential physics, therefore, lies not in the formation of the initial seed but in its subsequent resonant expansion and fluctuating interface growth described by KPZ universality.

Several observations indicate that the annealer probes genuinely quantum dynamics. First, the estimated time spent in the dynamical resonant regime is of order a few tens of nanoseconds, consistent with the timescales on which coherent behavior has previously been demonstrated on the same class of annealers [40–42]; see SI for the estimation. Second, the observed scaling $M(t) \sim h_x^2 t$, seen in both TTN calculations and quantum simulations, arises from resonance broadening, an intrinsically quantum effect. Third, we find that the classical-energy contribution remains approximately conserved over the experimentally relevant time window, showing that the observed growth stays concentrated near the resonant manifold rather than being dominated by generic off-resonant relaxation due to thermalization. These points, together with the agreement with approximate classical methods, provide evidence that the annealer realizes the targeted quantum many-body dynamics of the full model over the experimentally relevant time window. Its ability to simulate large systems is particularly important for extracting the late-time properties of bubble growth, where finite-size effects can otherwise obscure the asymptotic behavior.

At the same time, our results expose an intriguing puzzle. The microscopic dynamics generate substantial entanglement and exhibit signatures consistent with coherent quantum evolution, yet the asymptotic growth laws, i.e., nearly ballistic expansion and KPZ-type interface broadening, are already captured by a SC description of the effective resonant model. This suggests that the large-scale statistical properties of resonant bubble growth may be remarkably insensitive to many microscopic quantum details. Determining precisely how quantum effects influence the asymptotic dynamics, therefore, remains an important open question. Moreover, it is interesting to consider if and how the KPZ universality class changes in higher dimensions $d > 2$, where the filaments are growing $(d - 1)$ -dimensional surfaces. Addressing these questions will likely require new dynamical control capabilities of quantum annealers that are not yet publicly available [53], or significant advances in numerical methods for higher-dimensional quantum dynamics, where current classical approaches remain substantially more limited.

More broadly, the present study should be viewed as a controlled isolation of one particular post-nucleation growth channel, as part of a larger class of similar resonances that appear at other longitudinal field values and in arbitrary dimensions. Generic FVD process will, in general, involve a mixture of nucleation and multiple competing growth channels. This suggests that resonant accretion may provide a broader route to growth-dominated metastable decay whenever local interface moves become energetically favored. An important outlook is therefore to understand whether

an analog of this mechanism can arise in quantum field theory, where true-vacuum expansion is usually assumed to be pressure-driven rather than mediated by stochastic accretion. Establishing when such locally enhanced interface growth can compete with conventional bubble-wall motion would connect the resonant mechanism identified here to more generic FVD settings.

ACKNOWLEDGEMENTS

We would like to acknowledge the helpful discussions with Miha Nevemšek and Lorenzo Ubaldi. J.-Y.D. acknowledges funding from the European Union’s Horizon 2020 research and innovation programme under the Marie Skłodowska-Curie Grant Agreement No. 101034413. Z.P. acknowledges support by the Leverhulme Trust Research Leadership Award RL-2019-015 and EPSRC Grant UKRI3780. M.L. acknowledges support by the Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) under Germany’s Excellence Strategy – EXC-2111 – 390814868. L.P. acknowledges support by the Horizon Europe program HORIZON-CL4-2022-QUANTUM-02-SGA via the project 101113690 (PASQuanS2.1), and the computational resources of the INFN Padova HPC cluster, funded by the NextGenerationEU Terabit project on PNRR – Avviso n. 3264 ”per il Rafforzamento e creazione di Infrastrutture di Ricerca”, Missione 4, ”Istruzione e Ricerca”. K.M. acknowledges support from the project Jülich UNified Infrastructure for Quantum computing (JUNIQ) that has received funding from the German Federal Ministry of Education and Research (BMBF) and the Ministry of Culture and Science of the State of North Rhine-Westphalia. G.H. would like to acknowledge the financial support from ARIS, P1-0040 Nonequilibrium Quantum System Dynamics. G.H. and J.V. would like to acknowledge the financial support from ERC AdG HIMMS – Hidden metastable mesoscopic states in quantum materials. J.V. would like to acknowledge the financial support from ARIS, P1-0416 Physics of Quantum Technologies and J1-70063 Quantum Simulation of Non-equilibrium Phenomena on Quantum Devices. The authors gratefully acknowledge the Jülich Supercomputing Centre (<https://www.fzjuelich.de/ias/jsc>) for funding this project by providing computing time on the D-Wave Advantage™ System JUPSI through the Jülich UNified Infrastructure for Quantum computing (JUNIQ).

AUTHOR CONTRIBUTIONS STATEMENT

G.H. performed and analyzed the quantum simulations. L.P. performed and analyzed classical modeling using tensor networks. J.-Y.D. and M.L. conducted and analyzed classical modeling using stochastic circuits. J.V. conceptualized the study. All authors participated in the discussions of the results and writing of the manuscript.

COMPETING INTERESTS STATEMENT

The authors declare no competing interests.

-
- [1] S. Coleman, Fate of the false vacuum: Semiclassical theory, *Phys. Rev. D* **15**, 2929 (1977).
- [2] M. Voloshin, I. Y. Kobzarev, and L. B. Okun, On bubbles in metastable vacuum, *Yad. Fiz.* **20**, 1229 (1974).
- [3] S. Coleman and F. De Luccia, Gravitational effects on and of vacuum decay, *Phys. Rev. D* **21**, 3305 (1980).
- [4] A. Linde, Decay of the false vacuum at finite temperature, *Nuclear Physics B* **216**, 421 (1983).
- [5] S. Sachdev, Quantum phase transitions, *Physics world* **12**, 33 (1999).
- [6] F. Iachello and N. V. Zamfir, Quantum phase transitions in mesoscopic systems, *Phys. Rev. Lett.* **92**, 212501 (2004).
- [7] A. de la Torre, D. M. Kennes, M. Claassen, S. Gerber, J. W. McIver, and M. A. Sentef, Colloquium: Nonthermal pathways to ultrafast control in quantum materials, *Rev. Mod. Phys.* **93**, 041002 (2021).
- [8] P. Hänggi, P. Talkner, and M. Borkovec, Reaction-rate theory: fifty years after Kramers, *Rev. Mod. Phys.* **62**, 251 (1990).
- [9] A. J. Leggett, S. Chakravarty, A. T. Dorsey, M. P. A. Fisher, A. Garg, and W. Zwerger, Dynamics of the dissipative two-state system, *Rev. Mod. Phys.* **59**, 1 (1987).
- [10] K. Macieszczak, M. Gutta, I. Lesanovsky, and J. P. Garrahan, Towards a theory of metastability in open quantum dynamics, *Phys. Rev. Lett.* **116**, 240404 (2016).
- [11] C. Yin, F. M. Surace, and A. Lucas, Theory of metastable states in many-body quantum systems, *Phys. Rev. X* **15**, 011064 (2025).
- [12] A. H. Guth, Inflationary universe: A possible solution to the horizon and flatness problems, *Phys. Rev. D* **23**, 347 (1981).
- [13] G. Isidori, G. Ridolfi, and A. Strumia, On the metastability of the Standard Model vacuum, *Nuclear Physics B* **609**, 387 (2001).
- [14] G. Degrassi, S. Di Vita, J. Elias-Miro, J. R. Espinosa, G. F. Giudice, G. Isidori, and A. Strumia, Higgs mass and vacuum stability in the standard model at NNLO, *Journal of High Energy Physics* **2012**, 1 (2012).
- [15] C. Caprini, M. Chala, G. C. Dorsch, M. Hindmarsh, S. J. Huber, T. Konstandin, J. Kozaczuk, G. Nardini, J. M. No, K. Rummukainen, *et al.*, Detecting gravitational waves from cosmological phase transitions with LISA: an update, *Journal of Cosmology and Astroparticle Physics* **2020** (03), 024.
- [16] A. N. Kuznetsov and P. G. Tinyakov, False vacuum decay induced by particle collisions, *Phys. Rev. D* **56**, 1156 (1997).
- [17] L. Pavešić, M. Di Liberto, and S. Montangero, Scattering and induced false vacuum decay in the two-dimensional quantum Ising model, arXiv e-Prints (2026), [arXiv:2509.02702](https://arxiv.org/abs/2509.02702) [quant-ph].
- [18] R. Gregory, I. G. Moss, and B. Withers, Black holes as bubble nucleation sites, *Journal of High Energy Physics* **2014**, 1 (2014).
- [19] G. Gelmini and M. Gleiser, Kinetics of sub-critical bubbles and the electroweak transition, *Nuclear Physics B* **419**, 129 (1994).
- [20] K. Enqvist and I. Vilja, Induced nucleation in weak first order phase transitions, *Physics Letters B* **344**, 98 (1995).
- [21] A. H. Guth and E. J. Weinberg, Cosmological consequences of a first-order phase transition in the su_5 grand unified model, *Phys. Rev. D* **23**, 876 (1981).
- [22] S. Sarangi, G. Shiu, and B. Shlaer, Rapid tunneling and percolation in the landscape, *International Journal of Modern Physics A* **24**, 741 (2009).
- [23] M. Gleiser and R. C. Howell, Resonant nucleation, *Phys. Rev. Lett.* **94**, 151601 (2005).
- [24] X. Wang, A. Yosifov, A. Iyer, and J. Sun, Entanglement-facilitated macroscopic cluster formation in quantum many-body dynamics (2026), [arXiv:2605.22947](https://arxiv.org/abs/2605.22947) [quant-ph].
- [25] S. B. Rutkevich, Decay of the metastable phase in $d = 1$ and $d = 2$ Ising models, *Phys. Rev. B* **60**, 14525 (1999).
- [26] G. Lagnese, F. M. Surace, M. Kormos, and P. Calabrese, False vacuum decay in quantum spin chains, *Phys. Rev. B* **104**, L201106 (2021).
- [27] O. Pomponio, M. A. Werner, G. Zarand, and G. Takacs, Bloch oscillations and the lack of the decay of the false vacuum in a one-dimensional quantum spin chain, *SciPost Phys.* **12**, 061 (2022).
- [28] S. Darbha, M. Kornjača, F. Liu, J. Balewski, M. R. Hirsbrunner, P. L. S. Lopes, S.-T. Wang, R. Van Beeumen, D. Camps, and K. Klymko, False vacuum decay and nucleation dynamics in neutral atom systems, *Phys. Rev. B* **110**, 155103 (2024).
- [29] A. Zenesini, A. Berti, R. Cominotti, C. Rogora, I. G. Moss, T. P. Billam, I. Carusotto, G. Lamporesi, A. Recati, and G. Ferrari, False vacuum decay via bubble formation in ferromagnetic superfluids, *Nature Physics* **20**, 558 (2024).
- [30] R. Cominotti, C. Baroni, C. Rogora, D. Andreoni, G. Guarda, G. Lamporesi, G. Ferrari, and A. Zenesini, Observation of temperature effects on false vacuum decay in atomic quantum gases, *Phys. Rev. Lett.* **135**, 183401 (2025).
- [31] Y.-X. Chao, P. Ge, Z.-X. Hua, C. Jia, X. Wang, X. Liang, Z. Yue, R. Lu, M. K. Tey, X. Wang, and L. You, Probing false vacuum decay and bubble nucleation in a Rydberg atom array, *Phys. Rev. Lett.* **136**, 120407 (2026).
- [32] B. Song, S. Dutta, S. Bhawe, J.-C. Yu, E. Carter, N. Cooper, and U. Schneider, Realizing discontinuous quantum phase transitions in a strongly correlated driven optical lattice, *Nat. Phys.* **18**, 259 (2022).
- [33] Z.-H. Zhu, Y. Liu, G. Lagnese, F. M. Surace, W.-Y. Zhang, M.-G. He, J. C. Halimeh, M. Dalmonte, S. C. Morampudi, F. Wilczek, Z.-S. Yuan, and J.-W. Pan, Probing false vacuum decay on a cold-atom gauge-theory quantum simulator, arXiv e-Prints (2024), [arXiv:2411.12565](https://arxiv.org/abs/2411.12565) [cond-mat.quant-gas].
- [34] D. Luo, F. M. Surace, A. De, A. Lerose, E. R. Bennewitz, B. Ware, A. Schuckert, Z. Davoudi, A. V. Gorshkov, O. Katz, *et al.*, Quantum simulation of bubble nucleation across a quantum phase transition, arXiv e-Prints (2025), [arXiv:2505.09607](https://arxiv.org/abs/2505.09607) [quant-ph].
- [35] J. Vodeb, J.-Y. Desaulles, A. Hallam, A. Rava, G. Humar, D. Willsch, F. Jin, M. Willsch, K. Michielsen, and Z. Papić, Stirring the false vacuum via interacting quantized bubbles on a 5,564-qubit quantum annealer, *Nature physics* **21**, 386

- (2025).
- [36] G. Shavit, S. Nadj-Perge, and G. Refael, Ephemeral superconductivity atop the false vacuum, *Nature Communications* **16**, 2047 (2025).
- [37] U. Borla, A. Lazarides, C. Groß, and J. C. Halimeh, Microscopic dynamics of false vacuum decay in the 2 + 1D quantum Ising model, arXiv e-Prints (2026), [arXiv:2601.04305 \[quant-ph\]](#).
- [38] P. Osterholz, F. Bensch, S. Tang, S. B. Sheela, B. Sbierski, I. Lesanovsky, and C. Groß, Collective cluster nucleation dynamics in quantum magnets, arXiv e-Prints (2026), [arXiv:2512.04656 \[cond-mat.quant-gas\]](#).
- [39] See the Supplemental Online Material for further details of the experiments and numerical results.
- [40] A. D. King, S. Suzuki, J. Raymond, A. Zucca, T. Lanting, F. Altomare, A. J. Berkley, S. Ejtemaee, E. Hoskinson, S. Huang, E. Ladizinsky, A. J. R. MacDonald, G. Marsden, T. Oh, G. Poulin-Lamarre, M. Reis, C. Rich, Y. Sato, J. D. Whittaker, J. Yao, R. Harris, D. A. Lidar, H. Nishimori, and M. H. Amin, Coherent quantum annealing in a programmable 2,000 qubit ising chain, *Nature Physics* **18**, 1324 (2022).
- [41] A. D. King, J. Raymond, T. Lanting, R. Harris, A. Zucca, F. Altomare, A. J. Berkley, K. Boothby, S. Ejtemaee, C. Enderud, E. Hoskinson, S. Huang, E. Ladizinsky, A. J. R. MacDonald, G. Marsden, R. Molavi, T. Oh, G. Poulin-Lamarre, M. Reis, C. Rich, Y. Sato, N. Tsai, M. Volkmann, J. D. Whittaker, J. Yao, A. W. Sandvik, and M. H. Amin, Quantum critical dynamics in a 5,000-qubit programmable spin glass, *Nature* **617**, 61 (2023).
- [42] A. D. King, A. Nocera, M. M. Rams, J. Dziarmaga, R. Wiersema, W. Bernoudy, J. Raymond, N. Kaushal, N. Heinsdorf, R. Harris, K. Boothby, F. Altomare, M. Asad, A. J. Berkley, M. Boschnak, K. Chern, H. Christiani, S. Cibere, J. Connor, M. H. Dehn, R. Deshpande, S. Ejtemaee, P. Farre, K. Hamer, E. Hoskinson, S. Huang, M. W. Johnson, S. Kortas, E. Ladizinsky, T. Lanting, T. Lai, R. Li, A. J. R. MacDonald, G. Marsden, C. C. McGeoch, R. Molavi, T. Oh, R. Neufeld, M. Norouzpour, J. Pasvolsky, P. Poitras, G. Poulin-Lamarre, T. Prescott, M. Reis, C. Rich, M. Samani, B. Sheldan, A. Smirnov, E. Sterpka, B. T. Clavera, N. Tsai, M. Volkmann, A. M. Whitar, J. D. Whittaker, W. Wilkinson, J. Yao, T. J. Yi, A. W. Sandvik, G. Alvarez, R. G. Melko, J. Carrasquilla, M. Franz, and M. H. Amin, Beyond-classical computation in quantum simulation, *Science* **388**, 199 (2025).
- [43] P. Silvi, V. Giovannetti, S. Montangero, M. Rizzi, J. I. Cirac, and R. Fazio, Homogeneous binary trees as ground states of quantum critical Hamiltonians, *Phys. Rev. A* **81**, 062335 (2010).
- [44] J. Haegeman, C. Lubich, I. Oseledets, B. Vandereycken, and F. Verstraete, Unifying time evolution and optimization with matrix product states, *Phys. Rev. B* **94**, 165116 (2016).
- [45] D. Bauernfeind and M. Aichhorn, Time dependent variational principle for tree tensor networks, *SciPost Phys.* **8**, 024 (2020).
- [46] O. R. Bellwood, H. P. Casagrande, and W. J. Munro, Fractal path strategies for efficient two-dimensional density matrix renormalization group simulations, *Phys. Rev. B* **112**, 224433 (2025).
- [47] F. Balducci, A. Gambassi, A. Lerose, A. Scardicchio, and C. Vanoni, Localization and melting of interfaces in the two-dimensional quantum Ising model, *Phys. Rev. Lett.* **129**, 120601 (2022).
- [48] L. Pavešić, D. Jaschke, and S. Montangero, Constrained dynamics and confinement in the two-dimensional quantum Ising model, *Phys. Rev. B* **111**, L140305 (2025).
- [49] M. Žnidarič, Spin transport in a one-dimensional anisotropic Heisenberg model, *Phys. Rev. Lett.* **106**, 220601 (2011).
- [50] M. Ljubotina, M. Žnidarič, and T. Prosen, Spin diffusion from an inhomogeneous quench in an integrable system, *Nature Communications* **8**, 16117 (2017).
- [51] E. Ilievski, J. De Nardis, M. Medenjak, and T. Prosen, Superdiffusion in one-dimensional quantum lattice models, *Phys. Rev. Lett.* **121**, 230602 (2018).
- [52] M. Kardar, G. Parisi, and Y.-C. Zhang, Dynamic scaling of growing interfaces, *Phys. Rev. Lett.* **56**, 889 (1986).
- [53] R. Deshpande, M. Kheirkhah, C. Rich, R. Harris, J. Raymond, E. Hoskinson, P. Sathe, A. J. Berkley, S. Paul, B. Barch, D. A. Lidar, M. Müller, G. Aeppli, A. D. King, and M. H. Amin, Analog-digital quantum computing with quantum annealing processors, arXiv e-Prints (2026), [arXiv:2603.15534 \[quant-ph\]](#).
- [54] D-Wave Systems, *Technical Description of the D-Wave Quantum Processing Unit*, Tech. Rep. (D-Wave Systems Inc., Burnaby, BC, Canada, 2020) D-Wave User Manual 09-1109A-V.
- [55] S. Bravyi, D. P. DiVincenzo, and D. Loss, Schrieffer-wolff transformation for quantum many-body systems, *Annals of Physics* **326**, 2793 (2011).
- [56] F. Baccari, D. Bacilieri, M. Ballarin, F. P. Barone, F. Campaioli, A. G. Catalano, G. Cataldi, A. Coppi, A. Costantini, A. Datta, A. D. Girolamo, D. Jaschke, S. B. Kožić, G. Magnifico, C. Mordini, S. Montangero, S. Notarnicola, A. Pagano, L. Pavesic, D. Rattacaso, M. Rigobello, N. Reinić, S. Scarlatella, I. Siloi, P. Silvi, M. Tesoro, G. Torre, and D. Wanisch, *Quantum tea: qtealeaves* (2025).

METHODS

Quantum simulation on D-Wave's quantum annealer

Our quantum simulations utilized the D-Wave Advantage2_system1.6 and Advantage2_system1.13 quantum annealers, comprising 4593 and 4579 physical qubits, respectively. Of these, 4058 and 4044 physical qubits were used to implement the logical torus geometry shown in Fig. 1a,b. The embedding realizes a square lattice with periodic

boundary conditions in both directions, while bypassing missing qubits on the quantum processing unit.

The programmable D-Wave Hamiltonian takes the form [54]

$$\hat{\mathcal{H}}_{\text{DW}}(t) = -\frac{A(s)}{2} \sum_i \hat{\sigma}_i^x + \frac{B(s)}{2} \left(g(t) \sum_i h_i \hat{\sigma}_i^z + \sum_{i>j} J_{ij} \hat{\sigma}_i^z \hat{\sigma}_j^z \right), \quad (3)$$

where $s = t/t_a \in [0, 1]$ is the normalized anneal parameter, $A(s)$ sets the transverse-field energy scale, $B(s)$ the problem-Hamiltonian scale, and $g(t)$ is the time-dependent gain multiplying the linear biases h_i . Restricting J_{ij} to nearest-neighbor ferromagnetic couplings on the logical lattice, with $J_{\text{prog}} = -0.3$, and taking uniform biases $h_i = h_{\text{prog}}$, the Hamiltonian reduces, after factoring out $B(s)/2$, to the mixed-field Ising form in Eq. (1), with $h_x = A(s)/[B(s)J_{\text{prog}}]$ and $h_z(t) = -g(t)h_{\text{prog}}/J_{\text{prog}}$.

To realize the dynamics, we used a reverse-annealing protocol specified by the schedule $[(0, 1), (t_1, s_{h_x}), (t_2, s_{h_x}), (t_3, 1)]$, where s_{h_x} is chosen such that $h_x = A(s_{h_x})/(B(s_{h_x})J_{\text{prog}})$. Starting from a classical spin configuration at $h_x = 0$, the system is reverse annealed to a point with small but finite transverse field, held there for a pause time $\tau = t_2 - t_1$, and then annealed back to $h_x = 0$ for readout. Typical values are $t_1 = 1.0 \mu\text{s}$, τ ranging from $0.25 \mu\text{s}$ to $7.0 \mu\text{s}$, and $t_3 - t_2 = 200 \text{ ns}$.

We employed two protocols for the longitudinal field using the device's h -gain feature. In the first, $g(t)$ is held constant during the pause, corresponding to a constant $h_z(t)$; sweeping this value produces the resonance map shown in Fig. 1d. In the second, $g(t)$ is linearly ramped from $g(t_1) = 1$ to $g(t_2) = -1$, with $h_{\text{prog}} = -0.7$, so that $h_z(t)$ approaches the $|h_z| = 2J$ resonance while h_x remains finite. This isolates the fast interface-growth dynamics by ensuring that the system spends only a small fraction of the full reverse anneal near the resonance.

Effective Hamiltonian at the resonance

At first order in h_x/J , the dynamics at the resonance $|h_z| = 2J$ is governed by an effective constrained Hamiltonian,

$$\hat{\mathcal{H}}_{\text{eff}} = -h_x \sum_j \hat{\mathcal{P}}_j \hat{\sigma}_j^x, \quad \hat{\mathcal{P}}_j = \begin{cases} 1, & \text{if } \sum_{i \in \langle i, j \rangle} \sigma_i^z = 2, \\ 0, & \text{otherwise,} \end{cases} \quad (4)$$

where the projector $\hat{\mathcal{P}}_j$ enforces the local resonance condition: a spin may flip only if it has exactly one nearest neighbor in the true vacuum and three in the false vacuum. This projector structure makes the resonance highly selective: isolated nucleation events are strongly suppressed, while growth along an existing interface is allowed. The resulting dynamics is therefore not that of conventional pressure-driven bubble expansion, but a kinetically constrained spreading process in which the FVD is controlled by local moves at the boundary of preexisting bubbles.

The Hamiltonian (4) is obtained using a Schrieffer-Wolf transformation [55] with $\hat{H}_0$ given by the classical energy $-J \left(\sum_{\langle i, j \rangle} \hat{\sigma}_i^z \hat{\sigma}_j^z - 2 \sum_j \hat{\sigma}_j^z \right)$ and the perturbation $\hat{V}$ by the dynamical term $-h_x \sum_j \hat{\sigma}_j^x$. The classical energy $\hat{H}_0$ is conserved at all orders in this expansion; however, the bubble growth that we observe is more constrained than simply conserving this quantity. For example, the effective Hamiltonian (4) displays Hilbert space fragmentation, and the system only explores the sector of the initial state.

We are interested in sectors in which there can be multiple initial bubbles, but all of size 1 (i.e., corresponding to a single-site defect). In this case, there are two different emergent conserved quantities: the conservation of the number of bubbles and the absence of loops. The latter means that there are no sequences of neighboring (vertically or horizontally, but not diagonally) down-spins that form a closed path. For example,

$$\begin{bmatrix} \downarrow & \downarrow & \uparrow & \downarrow & \downarrow \\ \downarrow & \uparrow & \uparrow & \downarrow & \downarrow \\ \downarrow & \downarrow & \downarrow & \downarrow & \downarrow \end{bmatrix} \text{ and } \begin{bmatrix} \downarrow & \downarrow & \downarrow \\ \downarrow & \uparrow & \downarrow \\ \downarrow & \downarrow & \uparrow \end{bmatrix} \text{ are not loops, but } \begin{bmatrix} \downarrow & \downarrow & \downarrow \\ \downarrow & \uparrow & \downarrow \\ \downarrow & \downarrow & \downarrow \end{bmatrix}, \begin{bmatrix} \uparrow & \downarrow & \downarrow & \downarrow & \uparrow \\ \downarrow & \downarrow & \uparrow & \downarrow & \downarrow \\ \downarrow & \downarrow & \downarrow & \downarrow & \downarrow \end{bmatrix} \text{ and } \begin{bmatrix} \downarrow & \downarrow \\ \downarrow & \downarrow \end{bmatrix} \text{ are.} \quad (5)$$

For any finite value of h_x , the higher-order processes are still relevant, even though they are suppressed. Notably, already at the second order, the number of bubbles and the number of loops are no longer conserved. For example, the two processes below happen with a rate $-h_x^2/(8J)$. They both create a loop and split a single bubble into two as

$$\begin{bmatrix} \downarrow & \downarrow & \downarrow \\ \downarrow & \uparrow & \uparrow \\ \uparrow & \downarrow & \downarrow \end{bmatrix} \rightarrow \begin{bmatrix} \downarrow & \downarrow & \downarrow \\ \downarrow & \downarrow & \uparrow \\ \uparrow & \downarrow & \downarrow \end{bmatrix} \quad \text{and} \quad \begin{bmatrix} \downarrow & \downarrow & \downarrow \\ \downarrow & \uparrow & \uparrow \\ \uparrow & \downarrow & \downarrow \end{bmatrix} \rightarrow \begin{bmatrix} \downarrow & \downarrow & \downarrow \\ \downarrow & \downarrow & \uparrow \\ \uparrow & \uparrow & \downarrow \end{bmatrix}. \quad (6)$$

In fact, for the different system sizes and sectors we tested, we find that the combination of the first- and second-order Hamiltonians leads to almost all states with a given value of $\hat{H}_0$ being in the same sector. Thus, for a finite value of h_x used in the experiment, it is expected that configurations violating the constraints appear at later times. Due to the very strong effect of the second-order term on fragmentation, we can neglect the effect of terms of order 3 and beyond. We note that if we look at a region of false-vacuum (i.e., with only up-spins), the smallest bubbles that can nucleate resonantly are 4-bubbles in a square configuration. As these only occur at fourth order, these processes are highly suppressed and can be neglected if there is any bubble already present.

Tensor network simulations

In order to fully capture the quantum effects of the bubble growth, we perform simulations with tensor networks. These provide a controlled approximation to the evolution of the full quantum model. We use tree tensor networks (TTNs), where the tensors are arranged into a binary tree structure. The leaves of the tree contain the physical indices of the lattice sites, whilst the auxiliary tensors in the upper layers help to transfer correlations across distant sites. This is critical for simulating two-dimensional systems with a one-dimensional tensor network. In such cases, the lattice has to be mapped down to one dimension using a space-filling curve, which necessarily produces effective long-range interactions. Nonetheless, TTNs capture long-range correlations better than the standard matrix product states [43], and it has been recently shown that TTNs can simulate certain dynamical processes of relatively large spin systems in two dimensions [17, 37].

For the simulations of the bubble growth, we start with an initial product state where all spins are aligned down in the Z direction, and a single nucleation center pointing up. We evolve the state with the Ising Hamiltonian Eq. (1), with $J = 1$, $h_z = 2$ and varying h_x . This is akin to an instantaneous quench directly into the resonant point.

We simulate the time evolution with the time-dependent variational principle (TDVP) algorithm [44, 45]. We use the single-tensor update version of the algorithm, and pad the tensors of the initial product state to the maximal bond dimension. For TTNs, this is a much more efficient approach compared to the two-site TDVP; because the tensors in the TTN are of rank-3 with all legs at the maximal bond dimension χ , the computational complexity of the single-site TDVP scales with $\mathcal{O}(\chi^4)$, while the two-site TDVP scales as $\mathcal{O}(\chi^6)$.

We have checked that the results are converged in the size of the discretized timestep, and use $dt = 0.01/J$. All simulations were performed at a set of different bond dimensions, and we checked that the observables of interest are converged with increasing bond dimension. Despite the relatively small value of h_x , we find that the TTNs cannot accurately capture all features of the dynamics beyond $tJ \sim 5$. Whilst the average quantities (such as the total magnetization) are well converged, their spatial distributions are not. This is best seen in Fig. 3. The tensor network simulations were performed using the Quantum Tea package [56].

Classical stochastic circuit simulations

To simulate large systems, we turn to a classically simulable stochastic circuit inspired by the effective model (4). The basic circuit has a “brickwork structure”, meaning that for each time step, there are two intermediate steps during which the odd ($x + y$ odd) and even ($x + y$ even) sites are evolved, respectively. This is to ensure that the evolution does not violate the constraint. In each evolution step, all flippable sites (both up to down and down to up) with the correct parity are identified. Each of these is then flipped with a probability p or left unflipped with probability $1 - p$. We use $p = 0.1$, which keeps the dynamics stochastic and avoids the deterministic limit obtained for $p = 1$. We note that the results are insensitive to the exact value of p as long as $p \ll 1$. The only influence of p is a rescaling of time. We show this invariance with respect to p in the SI.

In order to make sure that what we observe is a generic feature of the constraint and not an artifact of our exact evolution scheme, we have also tested an alternative circuit. Instead of having a brickwork structure, each step is composed of a single evolution step in which both odd and even sites can be updated. We consider flips on all sites satisfying the constraint (still with probability p) and subsequently resolve “conflicts”, where flips on neighboring sites would lead to violations of the constraints. We have explored different ways of resolving these “conflicts” while maximizing the number of flips, and found that the alternate circuit also shows ballistic growth (with a slightly different velocity) in all cases. See SI for more details and results.

Radially averaged bubble size

In Fig. 3 we show the radially averaged bubble size r , extracted from TN and SC simulations. In both cases, it is

471 computed from the local magnetization as

$$r = \sqrt{2 \sum_{x,y} \langle n_{x,y} \rangle [(x - x_0)^2 + (y - y_0)^2]}, \text{ with } x_0 = \sum \langle n_{x,y} \rangle x, \ y_0 = \sum \langle n_{x,y} \rangle y, \quad (7)$$

472 where $n_{x,y} = (1 - \sigma_{x,y}^z)/2$ is equal to one if a site is in the true vacuum (and so part of the bubble) and to zero if it
 473 is in the false vacuum.

Supplementary Files

This is a list of supplementary files associated with this preprint. Click to download.

- [Sl.pdf](#)