

Controlling the glass transition through active fluctuating interactions

Emir Sezik^{1,*}, Henry Alston^{1,2,†} and Thibault Bertrand^{1,‡}¹*Department of Mathematics, Imperial College London, South Kensington, London SW7 2AZ, United Kingdom*²*Laboratoire de Physique, École Normale Supérieure, CNRS, PSL Université, Sorbonne Université, Université de Paris, 75005 Paris, France*

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Fluctuating pairwise interactions are understood to drive fluidlike states in dense biological systems. These states find a broad range of functionalities, such as directing growth during morphogenesis and forming aggregates with heightened mechanical response. However, a tractable model capturing the role of microscopic fluctuating interactions in these structural transitions is crucially lacking. Here, we study a p -spin model with fluctuating pairwise couplings (of strength D_a and persistence time t_a) as a schematic model for interaction-mediated fluidization. We find that while stronger fluctuations suppress the glass transition, more persistent fluctuations have the opposite effect. We identify the presence of an emergent fluctuation-dissipation relation at long times. We numerically extract the critical temperature $T_g(D_a, t_a)$ from a scaling relation near the transition, illustrating how microscopic fluctuations control the glass transition.

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The collective fate of many-body systems often hangs on the smallest fluctuations. In the restless world of dense biological matter, interactions never stand still. Instead, pairwise forces are generically subject to fluctuations. For instance, the dynamic turnover of the actin-myosin network controls the adhesion strength between neighboring cells [1,2], whereas stochastic type-IV pili mechanics generate intermittent attractive forces between bacteria [3,4]. Across embryonic tissues and cellular aggregates, these fluctuating forces lead to structural transitions between rigid and fluid states which are crucial for controlling tissue growth and remodeling during morphogenesis, and collective motion in health and disease [1–13]; they endow living matter with a remarkable ability to switch between solidlike and liquidlike states going across a glass transition [14–18].

Glassy systems are characterized by disordered (liquidlike) yet mechanically rigid (solidlike) structures [19–23]. At low temperatures, their relaxation times exceed any laboratory scale, leaving them unable to equilibrate with their environment [21,24,25]. This arises from frustration among numerous metastable states [26,27] and leads to aging and memory effects [28–38], dynamical heterogeneity [20,21,24,39–44], ergodicity breaking [29,31,45,46], and fluctuation-dissipation violations [26,41,47–49]. To address structural transitions in living systems, however, one must consider active matter. Indeed, at high densities, motile active matter shows phase

behavior analogous to, yet distinct from, passive glasses [50–65]. Early studies revealed that self-propulsion can shift and reshape the glass transition, either fluidizing dense suspensions or, conversely, promoting kinetic arrest depending on the propulsion mechanism and persistence [50,66–69]. For nonmotile active matter, the framework of kinetic theory has proved successful in describing the emergent structures observed in many-particle systems with active pairwise forces [4,70–72], but its applicability is ultimately limited at high particle density.

A quantitative theory describing the glass transition remains elusive [73]. Even in thermal systems, exact analytical results are largely confined to infinite-dimensional mean-field models, such as the hard-sphere glass transition [27,74–76]. More tractable descriptions arose from mode-coupling theory [49,77–79], derived *from first principles* using the Mori-Zwanzig projection formalism [80,81]. In active systems, mode-coupling theory successfully describes how self-propulsion alters relaxation pathways and collective dynamics, thus leading to modified glass transition lines [82–87]. Recent works further highlighted that dynamical heterogeneities in active glasses acquire features absent in equilibrium [88].

While mode-coupling theory has its own limitations [21], it naturally links structural glass transitions to spin glass theory. The p -spin model [89,90]—a generalization of the Sherrington-Kirkpatrick model [91] to higher-order interactions—provided the first mean-field framework capturing the phenomenology of glass formation [89,90]. Its dynamics for $p = 3$ map onto mean-field mode-coupling descriptions of supercooled liquids [92–94] making it a cornerstone of the random first-order transition theory [95] and a minimal, analytically tractable setting to study dynamical arrest, metastability, and the Gardner transition in thermal systems [47,79,96–100].

*Contact author: emir.sezik19@imperial.ac.uk†Contact author: henry.alston@phys.ens.fr‡Contact author: t.bertrand@imperial.ac.uk

Nonequilibrium extensions of the p -spin model have since been used to explore these phenomena in driven and active systems [50,100–104]. Adding external driving to each spin, as a proxy for the persistent motion of spins and a minimal source of self-propulsion-like activity, produces either partial or full suppression of the glass transition, depending on the driving timescale relative to glassy dynamics. For example, nonreciprocal interactions between the spins, whose timescale matches that of the glassy dynamics, have been shown to fully suppress the glass transition [105], whereas the transition survives upon the addition of active noise with fixed internal timescale. When the transition is present, the nonequilibrium transition still shows two-step decay and diverging equilibration time, as in the thermal p -spin model.

In this Letter, we show how active couplings reshape the glass transition in the p -spin model. These dynamical couplings are representative of fluctuating microscopic pairwise forces in structural glasses, a *bona fide* source of activity, requiring a finite rate of energy dissipation [106,107]. This yields an analytically tractable model of a nonequilibrium glass transition controlled by active fluctuating interactions. Though mean-field, this model qualitatively captures the effect of fluctuating interactions on the structural properties of living and active matter. Studying the mean-field equations allows us to gain intuition, building the foundation for the study of systems with fluctuating interactions where spatial dependence is incorporated. We analyze both extrinsic and intrinsic interaction fluctuations, characterized by strength D_a and persistence time t_a , showing that they can suppress the glass transition and fluidize the system. Furthermore, we demonstrate that they generate two distinct fluctuation-dissipation relations (FDRs) valid over different timescales. Altogether, our results reveal a genuinely nonequilibrium glass transition inimitable in thermal p -spin models.

Active p -spin model with fluctuating couplings. We consider an extension of the p -spin model supplemented by dynamic p -wise couplings in the mean-field limit. The p -wise coupling $\eta_{i_1 \dots i_p}(t)$, unlike the usual quenched coupling $J_{i_1 \dots i_p}$, fluctuates in time according to dynamics *independent* from the spins. These intrinsic fluctuations break detailed balance and drive the system away from equilibrium. A thermal p -spin model is defined by the following Hamiltonian:

$$\mathcal{H}[\{\sigma\}] = - \sum_{i_1 < \dots < i_p}^N J_{i_1 \dots i_p} \sigma_{i_1} \cdots \sigma_{i_p} + \mu \sum_{i_1=1}^N \sigma_{i_1}^2, \quad (1)$$

where $J_{i_1 \dots i_p}$ is a quenched variable drawn from a Gaussian distribution with zero mean and variance $p!/(2N^{p-1})$, and μ is a Lagrange multiplier enforcing the spherical condition $\sum_{i=1}^N \sigma_i^2 = N$. The dynamics of this model under the effect of p -wise couplings is governed by the following Langevin equation:

$$\dot{\sigma}_{i_1}(t) = - \frac{\partial \mathcal{H}[\{\sigma\}]}{\partial \sigma_{i_1}} + \sum_{i_2 < \dots < i_p} \eta_{i_1 \dots i_p}(t) \sigma_{i_2}(t) \cdots \sigma_{i_p}(t) + \xi_{i_1}(t), \quad (2)$$

where $\xi_i(t)$ is a Gaussian white noise with zero mean satisfying $\langle \xi_i(t) \xi_j(t') \rangle = 2T \delta_{ij} \delta(t - t')$, and $\eta_{i_1 \dots i_p}(t)$ is assumed to

follow an Ornstein-Uhlenbeck process:

$$t_a \dot{\eta}_{i_1 \dots i_p}(t) = -\eta_{i_1 \dots i_p}(t) + \zeta_{i_1 \dots i_p}(t), \quad (3)$$

with $\zeta_{i_1 \dots i_p}(t)$ a Gaussian white noise with zero mean satisfying $\langle \zeta_{i_1 \dots i_p}(t) \zeta_{j_1 \dots j_p}(t') \rangle = D_a p! / (2N^{p-1}) \delta(t - t') \prod_{k=1}^p \delta_{i_k, j_k}$. Physically, in this model, D_a corresponds to the energy imparted to the system due to active fluctuating couplings, and t_a , their persistence time, controls the amount of time over which this injection occurs.

Here, we assumed that the p -wise fluctuations are independent and local, however, this need not be the case. In [108], we also consider (i) the case of extrinsic fluctuations, where the relative strength of fluctuations is the same for all couplings, and we show that the resulting effective equations of motion have the same form for both of these models, and (ii) the case of a general process only requiring the $\eta_{i_1 \dots i_p}(t)$ to be drawn from a Gaussian distribution such that $\langle \eta_{i_1 \dots i_p}(t) \rangle = 0$ and $\langle \eta_{i_1 \dots i_p}(t) \eta_{j_1 \dots j_p}(t') \rangle = D_a p! / (2N^{p-1}) F(t - t') \prod_{k=1}^p \delta_{i_k, j_k}$, where $F(\tau)$ quantifies the persistence of the noise.

Dynamical description in the fluid phase. Starting from Eqs. (2) and (3), one can derive a closed set of equations of motion for the correlation function,

$$C(t, t') = \frac{1}{N} \sum_{i=1}^N \langle \sigma_i(t) \sigma_i(t') \rangle, \quad (4)$$

and the response function,

$$R(t, t') = \frac{1}{N} \sum_{i=1}^N \left\langle \frac{\delta \sigma_i(t)}{\delta \xi_i(t')} \right\rangle. \quad (5)$$

These two functions fully characterize the dynamics of the system [108].

For high enough temperature T , the system is in a fluid phase where we can assume that the dynamics are stationary for long enough waiting times t' , i.e., $C(t, t') = C(t - t')$ and $R(t, t') = R(t - t')$ as $t' \rightarrow \infty$. In this regime, the equations of motion are [108]

$$\frac{\partial C(\tau)}{\partial \tau} = -\mu C(\tau) + \int_0^\tau d\tau' \Sigma(\tau - \tau') C(\tau') \quad (6a)$$

$$+ \int_0^\infty d\tau' [\Sigma(\tau + \tau') C(\tau') + D(\tau + \tau') R(\tau')],$$

$$\frac{\partial R(\tau)}{\partial \tau} = -\mu R(\tau) + \int_0^\tau d\tau' \Sigma(\tau - \tau') R(\tau') + \delta(\tau), \quad (6b)$$

$$\mu = T + \int_0^\infty d\tau' [\Sigma(\tau') C(\tau') + D(\tau') R(\tau')], \quad (6c)$$

where

$$\Sigma(\tau) = \frac{p(p-1)}{2} C^{p-2}(\tau) R(\tau) [1 + D_a F(\tau)], \quad (7a)$$

$$D(\tau) = \frac{p}{2} C^{p-1}(\tau) [1 + D_a F(\tau)], \quad (7b)$$

$$F(\tau) = \frac{1}{2t_a} e^{-|\tau|/t_a}, \quad (7c)$$

and μ enforces the spherical condition $C(0) = 1$. It can be easily verified that by setting $D_a = 0$, we recover the equations of motion that describe the fluid phase of the thermal p -spin

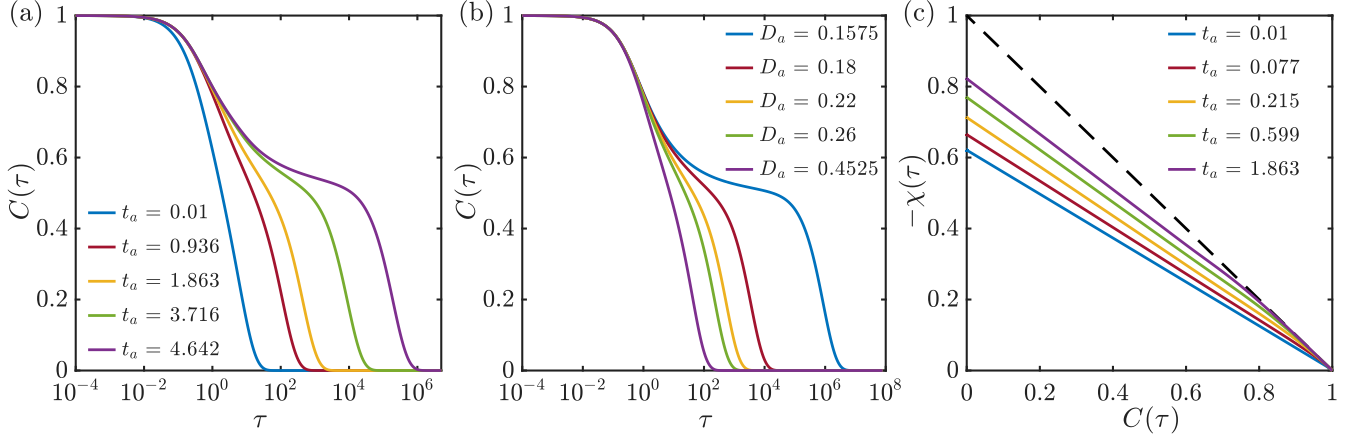


FIG. 1. Correlation functions and effective FDRs in fluid phase. (a) Correlation functions $C(\tau)$ for different values of the interaction fluctuation persistence time t_a , fixing $D_a = 0.4525$ and $T = 0.535$. (b) Correlation functions $C(\tau)$ shown for different interaction fluctuation strengths, D_a at $T = 0.555 < T_{g,\text{eq}}$ and $t_a = 0.6$. While increasing the interaction persistence time t_a drives the system towards the glass transition, increasing the fluctuation strength D_a has the opposite effect. (c) Integrated response function $\chi(\tau)$ as a function of the correlation function $C(\tau)$ for $T = 0.535$ and $D_a = 0.4525$. We identify two linear regimes, implying the existence of an effective FDR at both short and long times where the effective temperature is given by the gradient in each regime.

model [21,47,92,93]. In deriving the stationary equations, we have neglected the aging contribution, which dominates in the glass phase. Consequently, Eqs. (6) cease to hold within the glass phase, only capturing the onset of it.

Suppression of the glass transition by active fluctuations. Equations (6) can be fully understood analytically in the zero persistence limit, $t_a \rightarrow 0$, where fluctuations effectively become infinitely fast and $F(\tau) \rightarrow \delta(\tau)$. In this limit, the dynamics are that of a thermal p -spin model with effective temperature $T_{\text{eff}} = T + pD_a/4$ and the fluctuations purely heat up the system. Furthermore, because of this mapping to a thermal p -spin model, we also recover an *effective* FDR, $R(\tau) = -\frac{1}{T_{\text{eff}}} \frac{dC(\tau)}{d\tau}$. We use this to write a closed-form equation for $C(\tau)$:

$$\frac{dC(\tau)}{d\tau} = -T_{\text{eff}}C(\tau) - \frac{p}{2T_{\text{eff}}} \int_0^\tau d\tau' C^{p-1}(\tau - \tau') \frac{dC(\tau')}{d\tau'}, \quad (8)$$

which is precisely the equation of motion for a thermal p -spin model in the fluid phase. Standard results provide the glass transition temperature, T_g , which is present for $p \geq 3$ [47,92,93,108]:

$$T_g = \sqrt{\frac{p(p-2)^{p-2}}{2(p-1)^{p-1}}} - \frac{pD_a}{4}. \quad (9)$$

Notably, the glass transition is *completely suppressed* for strong enough coupling fluctuations,

$$D_a > D_a^* \equiv \sqrt{8(p-2)^{p-2}/[p(p-1)^{p-1}]}. \quad (10)$$

For $p = 3$, $D_a^* = \sqrt{\frac{2}{3}} \approx 0.81$.

In the glass phase, the correlation function $C(\tau)$ exhibits an infinite shoulder parametrized by the Edwards-Anderson parameter $C(\tau \rightarrow \infty) = q$. We determine q for the current model to solve [108]

$$T = \sqrt{\frac{p(p-1)q^{p-2}}{2}}(1-q) - \frac{pD_a}{4}. \quad (11)$$

We stress that, unlike the thermal p -spin model, q does not tend to 1 as $T \rightarrow 0$ as the active fluctuations always generate a residual temperature.

Infinitely persistent active fluctuations. In the limit $t_a \rightarrow \infty$, the fluctuations in the couplings disappear and the system reduces to the standard thermal p -spin model at temperature T . This reflects the fact that energy injected on timescales much longer than the relaxation time τ_{rel} is dynamically irrelevant: the fluid relaxes before the couplings evolve. In this regime, the activity effectively acts as an additional quenched disorder. However, it will still alter the behavior of the system due to the chaotic nature of the spin glasses [109–112]. Specifically, we expect the fluctuating couplings to distort the location of the ground-state energy as p -spin models have been shown to be chaotic under small perturbations to the disorder [113,114]. We expect this regime to hold as long as $t_a \gg \tau_{\text{rel}}$.

Active fluctuations with finite persistence. For finite t_a , the dynamics escape a direct mapping to that of a thermal p -spin model, thus describing a truly nonequilibrium model. Thus, we resort to solving Eq. (6) numerically (see [108] for details of the numerics). The stationarity assumption allows us to explore the correlation function up to $\tau = 10^7$ (i.e., over more than 10 decades).

Motivated by mode-coupling theory, we numerically analyze the $p = 3$ case, but we expect our results to be qualitatively valid for $p \geq 3$. Figures 1(a) and 1(b) show that fluctuating couplings stabilize the liquid phase well below the equilibrium glass temperature $T_{g,\text{eq}} = \sqrt{3/8}$ for finite activity strength D_a and persistence time t_a . Remarkably, D_a and t_a play opposite roles: increasing D_a fluidizes the system, while increasing t_a counteracts this effect. Physically, t_a cannot cool the system below T ; it offsets the heating induced by D_a .

The solutions $C(\tau)$ and $R(\tau)$ exhibit a clear timescale separation, with distinct effective FDRs for $\tau \ll t_a$ and $\tau \gg t_a$. Such generalized FDRs or effective temperatures have also proven to be essential in the description of the aging dynamics of thermal glassy systems [47,115]. Here, we use this to

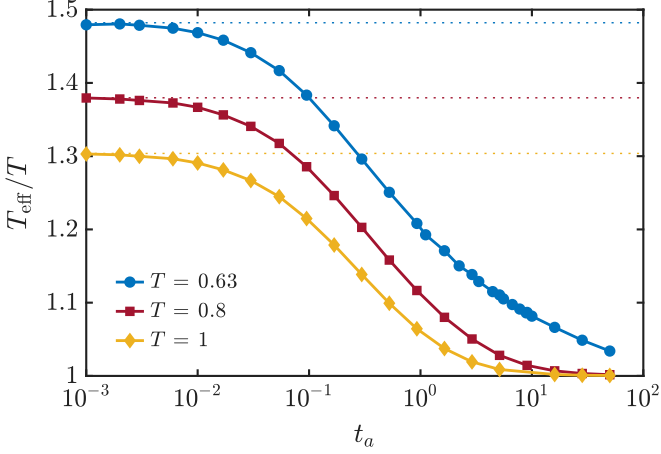


FIG. 2. Measured effective temperature for long-time FDR. Having identified the long-time FDR present in Fig. 1(c), we measure T_{eff} as a function of t_a for $D_a = 0.405$ and varied temperature T . The dashed lines denote the analytic result derived for T_{eff} in the white noise limit, i.e., a limit of nonpersistent fluctuations $t_a \rightarrow 0$.

quantify the timescale separation induced by the persistence of the activity t_a . Figure 1(c) shows the integrated response function $\chi(\tau) = -T \int_0^\tau d\tau' R(\tau')$ plotted against the correlation function. These curves display two linear regimes connected by a crossover, where the slopes define the effective temperatures. For $\tau \ll t_a$, one finds $T_{\text{eff}} = T$ as the fluctuations are effectively frozen, while for $\tau \gg t_a$, T_{eff} becomes nontrivial and depends on the activity parameters. The persistence time t_a determines the onset of this late-time regime: it occurs immediately for very small t_a , but is delayed for larger t_a .

In Fig. 2, we extract T_{eff} from the slope of Fig. 1(c). The results agree with the analytic limits found above for $t_a \rightarrow 0$ and $t_a \rightarrow \infty$. Convergence slows down near the equilibrium glass transition because lowering T increases τ_{rel} , requiring much larger t_a to ensure $t_a \gg \tau_{\text{rel}}$. Below $T_{g,\text{eq}}$, however, no such decay to T is observed: at these temperatures, the model

cannot be mapped onto a thermal p -spin system in the fluid phase even for $t_a \gg \tau_{\text{rel}}$, as it would already be in the glass phase. Thus, the thermal equations describing the fluid phase do not approximate our dynamics in this regime (see [108] for further discussion). The emergence of yet another distinct effective FDR in the aging regime might be expected as is typically seen in the thermal systems [47, 116, 117], but a numerical confirmation of this is beyond the scope of the current work.

The presence of these two effective FDRs and an internally fixed timescale t_a suggests a rewriting of the physical quantities in terms of “slow” and “fast” variables (e.g., $C = C_s + C_f$), as employed in Refs. [50, 102]. Integrating out the fast variables will result in an effective equation for the slow ones (identical to that in Ref. [50]):

$$\begin{aligned} \frac{\partial C_s}{\partial \tau} = & -(\mu - I_\Sigma)C_s(\tau) + I_R D_s(\tau) \\ & + \int_0^\infty d\tau' [\Sigma_s(\tau + \tau')C_s(\tau') + D_s(\tau + \tau')R_s(\tau')] \\ & + \int_0^\tau d\tau' \Sigma_s(\tau - \tau')C_s(\tau'), \end{aligned} \quad (12a)$$

$$\begin{aligned} \frac{\partial R_s}{\partial \tau} = & -(\mu - I_\Sigma)R_s(\tau) + I_R \Sigma_s(\tau) \\ & + \int_0^\tau d\tau' \Sigma_s(\tau - \tau')R_s(\tau'), \end{aligned} \quad (12b)$$

$$\begin{aligned} \mu = & T + \int_0^\infty d\tau' [\Sigma_f(\tau')C_f(\tau') + D_f(\tau')R_f(\tau')] \\ & + \int_0^\infty d\tau' [\Sigma_s(\tau')C_s(\tau') + D_s(\tau')R_s(\tau')], \end{aligned} \quad (12c)$$

where $I_\Sigma = \int_0^\infty d\tau \Sigma_f(\tau)$ and $I_R = \int_0^\infty d\tau R_f(\tau)$.

Combining the late-time effective FDRs with the slow dynamics yields a “closed” equation for $C_s(\tau)$, formally equivalent to a thermal p -spin model with activity renormalizing microscopic coefficients. This suggests that the glass transition survives and retains its essential features [50],

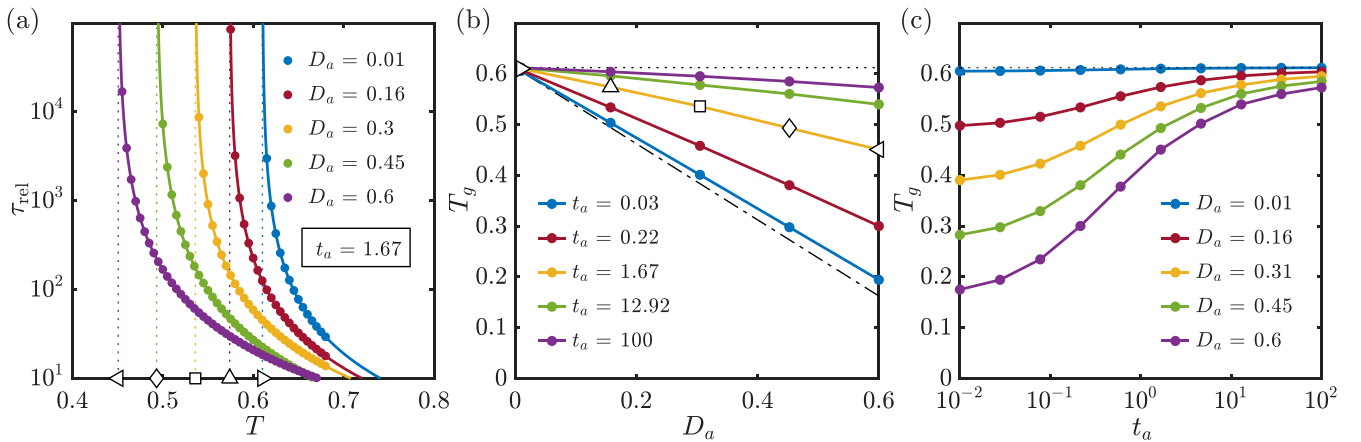


FIG. 3. Phase diagram for p -spin model with dynamical couplings. (a) The decay time τ_{rel} is plotted against T for $t_a = 1.67$ and varying D_a . We fit the data (dots) to a power law Eq. (14) (solid lines) to obtain the critical temperature T_g (vertical dotted lines) for each D_a, t_a . (b), (c) The resulting phase diagram for this model: as the strength (D_a) and persistence (t_a) of the coupling fluctuations vary, we evaluate $T_g(D_a, t_a)$ as in (a) and how it varies between the two fully tractable cases: (i) white noise fluctuations ($t_a \rightarrow 0$, dot-dashed line) and (ii) frozen couplings ($t_a \rightarrow \infty$, dotted line).

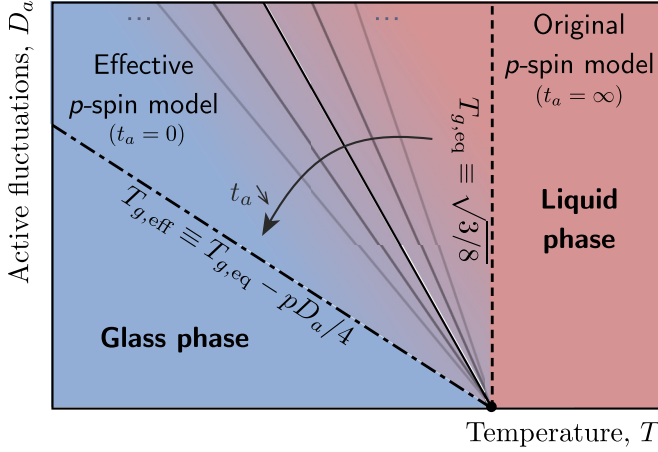


FIG. 4. Phase diagram. Schematic phase diagram of the p -spin model supplemented by active fluctuating couplings characterized by a strength D_a and a persistence time t_a . We show that the glass transition in this model can be controlled by two parameters (D_a, t_a); $T_g(D_a, t_a)$ varies smoothly between the original thermal p -spin model with $T_{g,eq}$ when $t_a \rightarrow \infty$ and an effective thermal p -spin model with $T_{g,eq} - pD_a/4$ when $t_a \rightarrow 0$. In particular, the glass transition can be partially suppressed by increasing D_a or lowering t_a . For large enough active fluctuations, complete suppression of the glass transition can be observed as $T_g(D_a, t_a)$ becomes negative.

which we confirm numerically. To probe it, we compute the relaxation time

$$\tau_{\text{rel}} = \int_0^\infty d\tau C(\tau). \quad (13)$$

In the thermal p -spin model, τ_{rel} diverges *algebraically* with exponent $\gamma = 1.765$ as we approach the glass transition [21]. As shown in Fig. 3, we find the same scaling form close to the critical temperature,

$$\tau_{\text{rel}} \sim (T - T_g)^{-\gamma}, \quad (14)$$

where the critical exponent γ and the glass transition temperature $T_g(D_a, t_a)$ are found through a fitting procedure [see Figs. 3(b) and 3(c)]. We find that γ is a nonuniversal exponent that depends explicitly on the activity parameters ranging between $\gamma \in [1.72, 3.1]$ (see [108] for a more detailed discussion).

Figure 4 succinctly summarizes our results. Increasing the activity strength D_a shifts the transition to lower temperatures and fluidizes the system, while larger persistence times t_a promote kinetic arrest. The critical temperature varies smoothly with t_a , interpolating between $T_{g,eq} - pD_a/4$ and $T_{g,eq}$; this dependence is shown explicitly in Fig. 3(c). At large D_a , the value of t_a becomes especially important in setting T_g .

Conclusions and outlook. We have demonstrated how dynamic couplings can partially or even fully suppress the glass transition temperature in a nonequilibrium p -spin model. We show that the strength and persistence of interaction fluctuations have opposite effects: stronger fluctuations reduce the glass transition temperature, whereas more persistent fluctuations alleviate this effect and induce an earlier onset of the glass phase. We demonstrate this by extracting the critical temperature from Eq. (14) as a function of the two parameters characterizing the fluctuations, $T_g(D_a, t_a)$.

Our results illustrate how fluctuating interactions drive transitions between glassy and fluid states. Interestingly, the tunability of the glass transition temperature T_g via the fluctuation strength D_a and persistence time t_a provides a possible microscopic mechanism for dense living systems to regulate how far from the glass transition they operate, enabling reversible switching between fluidlike and rigid states during morphogenetic or physiological processes. We expect that our framework can be extended to mixed spin models and models where the correlations between the fluctuating couplings follow a power-law decay (lacking any inherent timescale t_a), for which classification of the onset of glassy dynamics remains an intriguing open problem. More generally, we hope to extend this bottom-up approach, which directly connects microscopic dynamics to macroscopic transitions, to finite-dimensional systems exhibiting *structural* glass transitions. One route towards this would be to extend the mode-coupling theory for active matter to systems with fluctuating interaction forces.

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Data availability. The data that support the findings of this article are openly available [118].

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