A simple model for a non-Fermi liquid: The SYK Model

Georgios Chnitidis*
Department of Informatics, Technical University of Munich, Garching, 85748, Germany

Abstract This report consists a review of two published scientific papers regarding a simple model for a non-Fermi liquid: the SYK model. The articles that were reviewed, discuss the topic both from an experimental and theoretical point of view. The theoretical approach, is basically a proposal of an extended version of the Sachdev-Ye-Kitaev (SYK) model, for explaining a dynamical quantum phase transition from non-Fermi liquid fixed point to a Fermi liquid state with the property of deriving an exact solution for a respectively large N limit. The experimental one, incorporates a quantum simulation of the SYK model, since the physical realization of the SYK Hamiltonian is non-trivial and quite challenging, due to strong randomness and fully non-local fermion interactions. The model is simulated on a nuclear-spin-chain quantum simulator with main observations the fermion paring instability of the non-Fermi liquid and the chaotic - non-chaotic transition, and its results are promising for opening new horizons on the applications of the non-Fermi liquid states in quantum chaotic systems and the AdS/CFT duality.

I. INTRODUCTION

The Sachdev-Ye-Kitaev model [1, 2] is a strongly coupled, quantum many-body system that is chaotic, nearly conformally invariant, and exactly solvable. This remarkably unique combination of properties have driven the intense activity surrounding SYK and its applications within both high energy and condensed matter physics. As a quantum field theory, SYK constitute a new class of large N theories. The dominance of a simple and well-organized set of Feynman diagrams, iterations of melons, enables the computation of all correlation functions.

The SYK model is a more recent work on the initial SY model of many fermions with all-to-all interactions, a solvable example of a non-Fermi liquid with random couplings among strongly interacted Majorana fermions. The saddle point solution, is exact in the thermodynamic (large N) limit and realizes a non-trivial conformal fixed point with non-zero entropy density at vanishing temperature. The conducted studies showed a structure of $1/N$ fluctuations in the SYK model which indicated a connection to quantum gravity with AdS_2 black hole [1, 3–5], maximal chaotic behaviour for the generalized model with q-fermion interaction (where $q > 2$), derives a holographic duality to the $(1+1)d$ gravity with a bulk black hole. Hence, the model is of great interest to quantum information, string theory and condensed matter physics, as a solvable many-body system, which is the building block for constructing a metal, capturing some of the properties of non-Fermi liquids (NFL) [6], where the NFL is the "strange metal" phase at optimal doping of the cuprates high-temperature superconductors, where the resistivity scales linearly with temperature for a very large range in the phase diagram [4–6]. This generalized model is possible to drive a quantum phase transition between to low-energy fixed points which entails very different scrambling dynamics[7].

The questions in hand are, whether or not we can classify matter according to the way that it scrambles information, and if the SYK model is tunable with a classical phase transition to a different state without having the black hole scrambling levels [8]. Even though studying dynamical quantum phase transitions is essential for answering the aforementioned inquiries from a theoretical perspective, an experimental approach is of great significance too, since the derivation of an exact solution in strongly interacting chaotic quantum systems is extremely rare. The reasons for the rarity of experimental quantum simulations of the SYK model are due to the extreme computational power needed for simulating the Hamiltonian with strong randomness and fully non-local fermions interactions, and secondly the main Ansatz, i.e. the initial states of the simulated systems in terms of quantum states at different temperatures and measure the related dynamical properties. A high-fidelity, error-prone quantum simulator would be the solution to the prior problems.

II. THEORY: SOLVABLE MODEL FOR A DYNAMICAL QUANTUM PHASE TRANSITION FROM FAST TO SLOW SCRAMBLING

The generalization of the SYK model could provide some solid answers to the questions mentioned in the previous section. The generalization would enable a quantum phase transition between two low-energy fixed points, which involve different scrambling dynamics exactly computed in N large limit. The generalized SYK model described by the following equations and depict in Fig.1 [7]:

$$H = H_c + H_\psi + H_{c\psi} \tag{1}$$

with,

$$H_c = \frac{1}{(2N)^{3/2}} \sum_{ijkl} J_{ijkl} c_i^\dagger c_j^\dagger c_k c_l - \mu \sum_i c_i^\dagger c_i \tag{2}$$

* georgios.chitidis@tum.de

$$H_\psi = \frac{1}{M^{1/2}} \sum_{\alpha\beta} t_{\alpha\beta} \psi^\dagger_\alpha \psi_\beta - \mu \sum_\alpha \psi^\dagger_\alpha \psi_\alpha \tag{3}$$

$$H_{c\psi} = \frac{1}{(NM)^{1/4}} \sum_{i\alpha} \left(V_{i\alpha} c^\dagger_i \psi_\alpha + V^*_{i\alpha} \psi^\dagger_\alpha c_i \right) \tag{4}$$

with c fermions and properly anti-symmetrized J_{ijkl} being the random four-fermion coupling drawn from a Gaussian distribution with zero mean and $\overline{|J_{ijkl}|^2 = J^2}$. The ψ fermions are those on the peripheral sites $\alpha = 1, ..., M$.

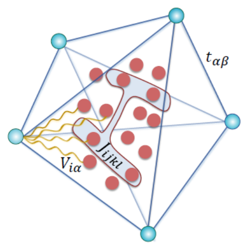

FIG. 1. The SYK sites at the centre are coupled through random four fermions coupling J_{ijkl}. The peripheral sites are connected to the SYK sites and to each other via random hopping $V_{i\alpha}$ and $t_{\alpha\beta}$ respectively [1].

For $N, M \to \infty$ proper thermodynamic limit with a fixed ratio $M/N = p$ is considered. The H_c Hamiltonian is solvable in the large-N limit and has an emergent conformal symmetry at low energies [2, 5, 7, 9, 10]. Thermalization and many-body quantum chaos with Lyapunov exponent $\lambda_L = 2\pi T (k_B = 1, \hbar = 1)$, are also arise from the generalized SYK model, saturating the quantum limit, like a black hole in Einstein gravity [2, 7, 10]. Regarding the ratio of peripheral and SYK sites p, if it is smaller than a critical value p_c then the dynamics are controlled by a strong coupling SYK-like fixed point with the universal Lyapunov exponent λ_L. If the ratio is bigger than a critical value the quadratic fermions screen the SYK interaction leading to a free low-energy fixed point with a non universal Lyapunov exponent [7]. This critical ratio p_c depends only on the fermion density and not the coupling strength, which sets the energy scale below the fermion coupling limit. Note that the proposed system will not always lead to a free fixed point in the low-energy limit, as a phase transition separating the two fixed point might be mistakenly assumed.

Furthermore, the Green's function for two species of the fermions is being used in order to derive the way that the quantum phase transition as a function p is manifested in the single-particle spectral properties. For $N \to \infty$ the Green's functions can be obtained either diagrammatically or, equivalently, from the saddle point of an effective action functional obtained via the replica formalism [7]. From here, the Green's function are being used to derive the Schwinger-Dyson equations:

$$G^{-1}(i\omega_n) = i\omega_n + \mu - \sum_J (i\omega_n) - V^2 \sqrt{p} \mathcal{G}(i\omega_n) \tag{5}$$

$$G^{-1}(i\omega_n) = i\omega_n + \mu - \sum_J (i\omega_n) - V^2 \sqrt{p} \mathcal{G}(i\omega_n) \tag{6}$$

with $\omega = (2n+1)\pi T$ the fermionic Matsubara frequency, n an integer and the last terms of both equations related to the self-energy diagrams. These equations are being used in order to derive the self energy at leading order in $1/N$, i.e.

$$\Sigma_J(\tau) = -J^2 G^2(\tau) \mathcal{G}(-\tau) \tag{7}$$

which is leading as well into two distinct phases at $T = 0$ tuned by the ratio p, the Non-Fermi Liquid phase (NFL) and the Fermi Liquid phase. Due to the conformal symmetry of the equations we get the following power-law forms:

$$G_R(\omega) = \Lambda \frac{e^{-i(\pi/4+\theta)}}{\sqrt{J\omega}} \tag{8}$$

$$\mathcal{G}_R(\omega) = \frac{\sqrt{p}}{V^2 \Lambda} e^{i(\pi/4+\theta)} \sqrt{J\omega} \tag{9}$$

$$\Sigma^R_J(\omega) = -\pi^{-1} \Lambda^3 e^{i(\pi/4+\theta)} \cos 2\theta \sqrt{J\omega} \tag{10}$$

with $\Lambda = \left(\frac{(1-p)\pi}{\cos 2\theta} \right)^{1/4}$, which determines the strength of the low frequency singularity, is having a singular behaviour during the tuning of the system towards the phase transition $p \to p_c$, while θ is related to spectral asymmetry and fermion filling through the Luttinger relation. The NFL state of the model is acquired for $p < p_c$, which is invariant under arbitrary re-parametrization of imaginary time $\tau = f(\sigma)$. Its characteristic is the singularity at $\omega \to 0$ in the single particle spectral functions. Two different behaviour for the density states of the two different species of the fermions, $c\psi$, which behaves as $1/\sqrt{\omega}$ and $\sqrt{\omega}$ respectively. The spectral functions vanish for $\omega \to 0^\pm$ on the boundaries, which might lead into a phase transition into an incompressible state. The numerical solution of the self consistency equations and the fermion density derivation as a function of the chemical potential at a fixed value of p, can help on the prediction of its occurrence. The calculation showed a direct transition to a compressible metallic Fermi liquid in values of p closer to 1 and an incompressible NFL state for smaller values of p. Whereas for a Fermi liquid the rate of the decay for the $T = 0$ quasi-particle vanishes as $\omega \to 0$. It is obtained that ω^2 diverges as $p \to p_c = 1$, which leads to the breakdown of the Fermi-liquid at the critical point. The saddle-point equations (5), (6), (7) have a trivial emergent conformal symmetry with scaling dimensions corresponding to non-interacting fermions [7]. The numerical results of the spectral function across the quantum critical points showed that in the NFL phase, they match at low-energies with the finite-temperature spectral densities obtained in the conformal limit, while after the transition to Fermi liquid phase the singularities are rounded off (even for T=0).

Due to the fact that the SYK model is characterized by a residual entropy, zero-temperature, for the case of the consideration of the thermodynamic limit $N \to \infty$ before the zero temperature limit, the low-energy spectrum of the SYK Hamiltonian is similar to the spectrum of usual quantum many-body high energy systems. The residual entropy is derived via the Maxwell relation and using the Luttinger relation we get:

$$S_0(n) = S_0(n_0) + \int_{n_0}^{n} dn \ln(\tan(\pi/4 + \theta(n))) \qquad (11)$$

which at fixed $p < 1$ the entropy vanishes. As a remark, the vanishing of the entropy at the critical points, possibly leads to a change of the geometry in the dual gravity formalism of the transition which means elimination of the black hole since the NFL to Fermi liquid critical point creates a chaotic quantum many-body transition.

The aforementioned conclusion is being elaborated in a more solid manner with the characterization of the quantum scrambling dynamics. This is achieved by means of the Lyapunov exponent λ_L for the four point out-of-time-ordered (OTO) correlation function, which leads to a saturating chaos [10] bound in the conformal low temperature limit over the entire NFL phase [7]. The Majorana fermions Hamiltonian that was used for the computing of the OTO is the following:

$$H = \frac{1}{4!}\sum_{ijkl} J_{ijkl}\chi_i\chi_j\chi_k\chi_l + \frac{i}{2!}\sum_{\alpha\beta} t_{\alpha\beta}\eta_\alpha\eta_\beta + i\sum V_{i\alpha}\chi_i\eta_\alpha \qquad (12)$$

with χ_i and η_i instead of c_i and ψ_i, the fully antisymmetric $t_{\alpha\beta}$ and J_{ijkl}. The OTO correlations for two SYK sites, which can help us prove quantum chaos and involves four Majorana operators is of the form [7]:

$$\langle\overline{\chi_i(t)\chi_j(0)\chi_i(t)\chi_j(0)}\rangle \simeq f_0 - \frac{f_1}{N}e^{\lambda_L t} + \mathcal{O}(\frac{1}{N^2}) \qquad (13)$$

and the time scale of its duration $t \lesssim t^* \simeq (1/\lambda_L)\ln(N)$, i.e. scrambling time, the decay time of the OTO correlation to small values and information encoded in local observables is lost to operators encompassing the entire system [10], with scrambling rate the Lyapunov exponent $\lambda_L \leq 2\pi T$. In the examined model, the coupling between the peripheral sites and the SYK fermions is such that the OTO correlation and correlation function are involved in a "cross scrambling" occurrence, thus having two coupled four-point functions $F_{\chi\chi\chi\chi} = F_1$ and $F_{\eta\eta\chi\chi} = F_2$ [7].

By solving the self consistent equations presented in the above figure, one can obtain the anticipating chaotic dynamics Ansatz for the function $\mathcal{F}$ in the form of an eigenvalue equation. Solving analytically the eigenvalue problem in the conformal limit $T \to 0$ in the NFL phase.

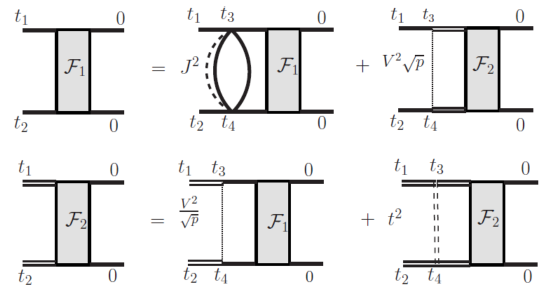

FIG. 2. The self-consistent equations for $1/N$ part of the OTO correlation function for Majorana fermions. Solid lines are the Green's function G and double lines $\mathcal{G}$ [7].

Finally, the single integral equation obtained is:

$$\frac{3}{4\pi}\frac{|\Gamma(\frac{1}{4}+\frac{h}{2}+iu)|^2}{|\Gamma(\frac{3}{4}+\frac{h}{2}+iu)|^2}\int_{-\infty}^{\infty} du'|\Gamma(\frac{1}{2}+i(u-u'))|^2 f_1(u') = (k-\frac{p}{k})f_1(u) \qquad (14)$$

One can acquire the solution of the self-consistent equations by setting $k = 1$ in $\frac{3(1-p)}{1+2h} = (k - \frac{p}{k})$ which leads to $h = 1$ and the Lyapunov exponent being equal to $\lambda_L = 2\pi T$. The conclusion is being made that the NFL phase saturates the chaos bound for $p < 1$ at half filling, a result which cannot be accounted at the QCP since the cut-off for the conformal limit vanishes at $p = 1$.

For arbitrary filling in the NFL phase of our complex fermion model, the computations of the low-temperature Lyapunov exponent are generalized. For the four-point functions, the depicted in figure 3 self-consistent approximation were used. Now by diagonalizing the single integral equation of the latter case by the same eigenfunction that solves the single integral equation (14), we get the

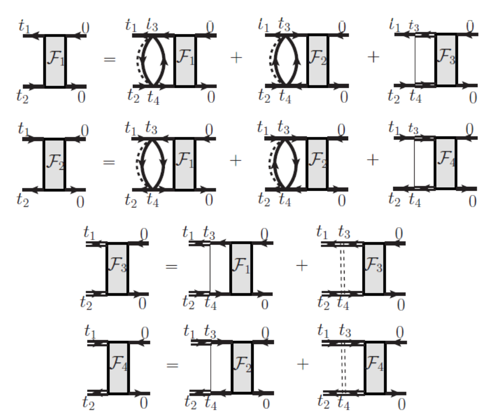

FIG. 3. The self-consistent approximations for the $1/N$ parts of the OTO correlation functions for complex fermions [7].

following eigenvalue equation:

$$\begin{pmatrix} 1 & \frac{1}{2}\tan(\frac{\pi}{4}+\theta) \\ \frac{1}{2}\cot(\frac{\pi}{4}+\theta) & 1 \end{pmatrix} \begin{pmatrix} a \\ b \end{pmatrix} = l_k(h) \begin{pmatrix} a \\ b \end{pmatrix} \quad (15)$$

and the eigenvalues were found to be, $l_k(h) = 1/2, 3/2$. The solution of the self consistent equation ($k = 1$) leads to $h = 0, 1$ which provides scrambling with Lyapunov exponent $\lambda_L = 2\pi T$ which yet again proves that the NFL phase saturates completely the chaos bound.

Finally, the acquired numerical results, i.e the computed Lyapunov exponent for half filling at finite temperature by means of the eigenvalue equation after applying the frequency ω discretization, indicate that the eigenvalue $k = 1$ is non-degenerate and data consistency with ratio of $h = \lambda_L \beta / 2\pi$ approaching to 1 in the NFL phase. Also for $p > 1$ there is consistency in the temperature dependency of the Lyapunov exponent with T^2 behaviour, which is expected for a Fermi liquid phase ($p \geq 2$).

III. EXPERIMENT: QUANTUM SIMULATION OF THE NON-FERMI-LIQUID STATE OF SACHDEV-YE-KITAEV MODEL

The goal of the experimental that is conducted in [8] is to investigate the fermion-pair instability of the SYK NFL and the recently predicted by [11] chaotic - non-chaotic transition. A 4-qubit NMR quantum simulator was used to simulate the $(0+1)d$ generalized SYK model with $N = 8$ Majorana fermions. With including measurement of the boson correlation function for different temperature and perturbations. Two different phases of the model exists: maximal chaotic NFL phase and perturbatively weak chaotic fermion-pair condensate phase. The Hamiltonian of the quantum system is given by

$$H = \frac{J_{ijkl}}{4!} X_i X_j X_k X_l + \frac{\mu}{4} C_{ij} C_{kl} X_i X_j X_k X_l \quad (16)$$

where $X_{i,j,k,l}$ the Majorana fermions and both J_{ijkl}, C_{ij} antisymmetric random tensors drawn from a Gaussian distribution. The pure SYK model is described for the Hamiltonian at $\mu = 0$ with low-temperature state in the limit $N \gg J/T \gg 0$ is maximally chaotic NFL. For $\mu > 0$ we have a marginally relevant perturbation in the Hamiltonian that leads to a spontaneous time-reversal symmetry breaking ($\mathcal{T} : X_i \to X_i, i \to -i$). The bosonic fermion-pair operator that is an outcome of the $\mathcal{T}$-breaking phase creates a persistent correlation $\langle b(t) b(0) \rangle$ that doesn't decay in time, and it can be used as its experimental signature. For a negative μ the spontaneous symmetry breaking [11] doesn't occur, thus the system is in maximally chaotic NFL phase. For the physical realization of the system the encoded Hamiltonian into the 1/2-spin operators via the Jordan-Wigner transformation yields:

$$\hat{H}_{\mathrm{NMR}} = \sum_{i=1}^{4} \frac{\omega_i}{2} \hat{\sigma}_z^i + \sum_{i<j,=1}^{4} \frac{\pi J_{ij}}{2} \hat{\sigma}_z^i \hat{\sigma}_z^j \quad (17)$$

with σ being the Pauli matrices, and ω_i the chemical shift of spin i and J_{ij} the coupling constant between spins. In FIG.1 the simulated quantum circuit is displayed. In or-

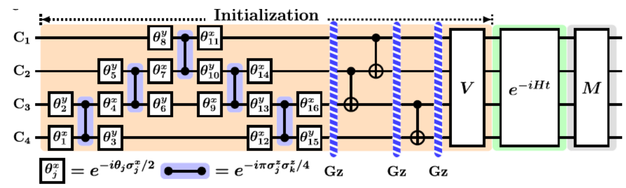

FIG. 4. The quantum circuit for measuring the boson correlation function with V the basis transformation of initial states ρ_i from the computational basis to the eigenvectors. Where M is the five readout pulses for observing the bosonic fermion-pair operator b.

der for the initial states to be prepared, the fact that the natural system is initially under the thermal equilibrium state $\rho_{eq} \approx (\mathbb{1} + \epsilon \sum_{i=1}^{4} \sigma_i^z)/2^4$, where $\epsilon \sim 10^{-5}$ is the polarization. Since there is not loss of unitarity we can omit the unit operator. We get the states with diagonal elements the eigenvalues of ρ_i. After the CNOT gates applied, the coherence of the quantum system was removed and remain unaffected by the z-direction gradient fields, after which the states are the diagonal density matrices, $\rho_i^d = V^\dagger \rho_i V$.

The evolution of generalized SYK model can be simulated with a controllable NMR system efficiently, as pointed out originally by Feynman [12–14]. The Hamiltonian (1) can be expressed as the sum of the spin interactions which yields:

$$H = \sum_{s=1}^{70} H_s = \sum_{s=1}^{70} a^s_{ijkl} \sigma^1_{a_i} \sigma^2_{a_j} \sigma^3_{a_k} \sigma^4_{a_l} \quad (18)$$

with $a = \{0, x, y, z\}$ being the labelling to the Pauli matrices. To incorporate a higher fidelity factor the use of Trotter-Suzuki formula is being used. Its exact time evolution operator can be decomposed into:

$$e^{-iH\tau} = \left(\prod_{s=1}^{70} e^{-iH_s \tau/n} \right)^n + \sum_{s<s'} \frac{[H_s, H_{s'}]\tau^2}{2n} + \mathcal{O}(|a|^3 \tau^3/n^2) \quad (19)$$

The accuracy of the approximation is being improved for large n and $[H_s, H_{s'}] = 0$. For the creation of local many-body spin interaction (H_s), the need for creating coherent control is emerging, which is satisfied by means of pulse sequence for finding the appropriate amplitudes and phases of radio-frequency fields. The algorithm used is called GRAPE (gradient ascent pulse engineering) for the field parameters of a shaped pulse optimization. An important indication that the quantum simulation algorithm is efficient is the fact that the total numbers of gates being used in the circuit grows polynomially with the number of Majorana fermions N.

The measurement of the boson correlation function is being made for examine whether or not the simulated

SYK NFL ground state is stable towards any spontaneous symmetry braking. The correlation function reads:

$$\langle b(\tau)b(0)\rangle_\beta = \frac{\mathrm{Tr}(e^{-\beta H}e^{-iH\tau}be^{iH\tau}b)}{\mathrm{Tr}(e^{-\beta H})} \tag{20}$$

Averaging the normalized correlation function over random samples is resolving the initial value fluctuation over the various samples. The normalization is being used for the exclusion of nonsensical phase interference among different samples. The measured NMR signal via quadrature detection reads:

$$S(t) = \mathrm{Tr}\left[\rho_f M^\dagger e^{iH_{\mathrm{NMR}}t}\sum_{j=1}^{4}(\sigma_x^j + i\sigma_y^j)e^{-iH_{\mathrm{NMR}}t}M\right] \tag{21}$$

where ρ_f is the output density matrix. NMR signal is comprised by both real and imaginary parts and is the average of transverse magnetization without any readout pulse. Here, five readout pulses are being used to get the $\mathrm{Tr}[\rho_f b]$ completely and consequently help us acquire the bosonic fermion-pair operator b.

The experimental results obtained by averaging over eight random samples with the truncation error due to the fitting procedure and fluctuations. The effect of decoherence was partially removed by means of normalization of the correlation function.

For a infinitely high temperature $\beta = 0$, the boson correlation functions seems to be invariant under the examined transformation $H \to -H$. However the correlation is decaying faster and to the lowest saturation value at $\mu = 0$, which was expected since the pure SYK model is maximally chaotic scrambling the order parameter.

For temperature corresponding to $\beta = 1$, the thermal fluctuation is destroying the boson condensation and the long-time correlation is suppressed.

Finally, for the low-temperature $\beta = 20$, the bosonic condensation is decaying quickly in the same pattern, i.e. chaotic phases. At $\mu = 0$ the pure SYK model is maximally chaotic NFL phase, whereas for $\mu = 5$ the decay is much slower and the saturation is happening to a relatively large value. The long-time order differences is actually the spontaneous $\mathcal{T}$-breaking phase and for $\mu < 0$ the systems is stable towards the symmetry breaking. Thus, the sign of μ is the catalytic factor on the continuity of the chaotic-non-chaotic transition, with critical properties analogous to the Kosterlitz-Thouless transition [8].

IV. CONCLUSIONS

Both papers [7, 8] examine the behaviour of a generalized SYK model and the many-body quantum chaos localization transition, with Majorana fermions residing in the core sites coupled to the ones on the peripheral sites (non-interacting). The ratio of the number of sites in the model is important for the quantum phase transition from the NFL phase in the pure SYK model to a Fermi liquid phase, with the residual entropy vanishing after passing the critical point. Different qualitative chaotic dynamical behaviour characterized by OTO correlations, obtained for the two different phases. The scrambling rate (i.e. Lyapunov exponent) which is the catalytic factor for defining the emergence of chaos, saturates the quantum bound in low-temperature limit ($\lambda_L = 2\pi T$) for the NFL phase, whereas for the Fermi liquid phase it is perturbative in the fermions interactions in the core sites ($\lambda_L \propto T^2$). It is proposed for the critical point, where the Lyapunov exponent is cancelled on either site of the transition, to be the new dynamical universality class and the holographic interpretation of the transition to be studied in more depth [7]. In the NFL phase the fast scrambling is mainly attributed to the AdS_2 spacetime correspondence emergence with a quantum black hole, which essential affects our perspective about the transition to the free fixed point and the fundamental geometry alternations that would eliminate the black hole. Regarding the quantum critical point, based on what was mentioned above, one can see that in the described model plays the role of the separator of the states with fast and slow scrambling. An interesting future approach would be any modifications of the proposed model towards the direction of a complete halt in the chaotic modes, while maintaining the analytical control in the N-large limit.

In the meanwhile, from an experimental point of view, are required better techniques for constructing high fidelity quantum circuits and algorithms in order to simulate the SYK model and any possible generalization. The NMR quantum simulator that is mentioned in part III, provides the strategic advantage [8] of well characterized qubits with long decoherence time and fine control of nuclear spins by means of radio-frequency fields enabling the simulating dynamics of the model. The maximal chaotic NFL phase's instability against specific types of four-fermions perturbations, is being brought up by the measurements of the fermion-pair correlation functions. This leads the model into a less chaotic condensed phase with spontaneous $\mathcal{T}$-breaking. The examined method can be combined with the generalized model that was described in part II for a deeper exploration of the holographic duality regime.Furthermore, it would be undoubtedly interesting to test non-equilibrium dynamics and the quantum information scrambling by measuring the Lyapunov exponent and the OTO correlation in order to define the chaos in the quantum system. The most challenging part in such an endeavour would be to create a quantum simulator with scalable control techniques and number of qubits, since the currently used gradient-based optimal control methods pose an obstacle towards that direction. To that end, a great amount of machine learning techniques are published every other day, in order to boost the efficiency of the quantum algorithms and enhance quantum control of the circuit, that would require polynomial time for optimization with the number of qubits. Such a simulator would be a handy tool if we were to simulate one of the most intrigued aspects of a more complex SYK model, which is a black hole and the experimental approach of the quantum gravity theory in a laboratory.

REFERENCES

[1] Subir Sachdev, Jinwu Ye *Gapless Spin-Fluid Ground State in a Random Quantum Heisenberg Magnet*, Phys. Rev. Lett. 70, 3339-3342 (1993). http://link.aps.org/doi/10.1103/PhysRevLett.70.3339.

[2] A. Kitaev, *A simple model of quantum holography*, KITP strings seminar and Entanglement 2015 program (Feb. 12, April 7, and May 27, 2015). http://online.kitp.ucsb.edu/online/entangled15/.

[3] Joseph Polchinski and Vladimir Rosenhaus, *The spectrum in the Sachdev-Ye-Kitaev model*, Journal of High Energy Physics 2016, 1-25 (2016).

[4] Subir Sachdev, *Holographic Metals and the Fractionalized Fermi Liquid*, Phys. Rev. Lett. 105, 151602 (2010).

[5] Juan Maldacena and Douglas Stanford, *Comments on the Sachdev-Ye-Kitaev model* (2016), arXiv:1604.07818.

[6] Vladimir Rosenhaus, *An introduction to the SYK model*, J. Phys. A: Math. Theor. 52 (2019) 323001, arXiv:1807.03334.

[7] Sumilan Banerjee, Ehud Altman, *Solvable model for a dynamical quantum phase transition from fast to slow scrambling*, Phys. Rev. B 95, 134302 (2017), arXiv:1610.04619.

[8] Zhihuang Luo, Yi-Zhuang You, Jun Li, Chao-Ming Jian, Dawei Lu, Cenke Xu, Bei Zeng, Raymond Laflamme, *Quantum Simulation of the Non-Fermi-Liquid State of Sachdev-Ye-Kitaev Model*, NPJ Quantum Information, 5, 7 (2019), arXiv:1712.06458.

[9] Subir Sachdev, *Bekenstein-hawking entropy and strange metals*, Physical Review X 5, 1-13 (2015), 1506.05111.

[10] Juan Maldacena, Stephen H. Shenker, and Douglas Stanford, *A bound on chaos*, Journal of High Energy Physics 2016, 1-17 (2016).

[11] Bi, Z., Jian, C. -M., You, Y. -Z., Pawlak, K. A. & Xu, C. *Instability of the non-fermi liquid state of the Sachdev-Ye-Kitaev model*, Phys. Rev. B 95, 205105, (2017).

[12] Gross, D. J. & Rosenhaus, V. *A generalization of Sachdev-Ye-Kitaev*, J. High. Energy Phys. 2, 93 (2017).

[13] Garca-Álvarez, L. et al., *Digital quantum simulation of minimal ads/cft*, Phys. Rev. Lett. 119, 040501 (2017).

[14] Feynman, R. P., *Simulating physics with computers*, Int. J. Theor. Phys. 21 467–488, (1982).