

# A Review on Quantum Operator Scrambling within Classically Simulable Circuit Models

Samuel A. Hopkins<sup>1</sup>

<sup>1</sup>*H. H. Wills Physics Laboratory, University of Bristol, Bristol, BS8 1TL, UK*

(Dated: November 24, 2022)

Abstract Goes Here...

## CONTENTS

I. INTRODUCTION	1
II. Background Theory	2
A. Quantum Information and Computing	2
III. Classically Simulable Quantum Circuits	3
A. Clifford Circuits and Stabilizer Formalism	3
B. Non-interacting Fermi Circuits	4
1. Fermionic Fock Space	4
2. Fermionic Circuits	5
IV. Operator Scrambling	5
A. Operator Spreading	5
B. Entanglement Growth in Operator Space	6
1. Blake and Linden's Construction	6
2. Scrambling in Free-Fermi Circuits	6
V. Further Direction	6
References	8

## I. INTRODUCTION

Understanding the evolution of quantum many-body systems presents challenging problems and has provided insightful results in condensed matter physics and quantum information [1]. Studying the dynamics of such systems is fundamentally a computational challenge due to the exponential growth of the Hilbert space with the number of qubits. However, quantum dynamics can be efficiently simulated in atypical quantum circuit models, which will be studied in hopes of understanding the spreading and scrambling of encoded local information, a process known as quantum scrambling. More precisely, quantum scrambling describes the process in which local information encoded by a simple product operator, becomes 'scrambled' by a unitary time evolution amongst the large number of degrees of freedom in the system, such that, the operator becomes a highly complicated sum of product operators with no decoding protocol to recover the information. In recent years, the study of

this scrambling process has yielded insightful results and new perspectives. These include new insights into black hole information and the AdS/CFT correspondence [2–8], quantum chaos, [9] and hydrodynamics in many-body systems [10–14].

In order to investigate this phenomenon, we turn to circuit models that are analytically and numerical tractable, so that a deeper understanding can be reached in hopes of applying any findings to generic cases of unitary evolution in many-body systems. As previously mentioned, some families of circuits are able to be efficiently simulated with polynomial effort on a classical computer. Two well-studied circuit models take the principal interest of this review, namely Clifford circuits and non-interacting fermi circuits. Clifford circuits and the Clifford group have played a key role in the study fault-tolerant quantum computing and, more recently, many-body physics via random Clifford unitaries [15, 16]. This is due to their efficient simulability on classical computers, with no constraints on the amount of entanglement, allowing for the dynamics of large numbers of qubits to be studied via stabilizer states [17].

Valiant [18] presented a new class of quantum circuits that can be simulated in polynomial time. This class of circuits was later identified and mapped onto a non-interacting fermion system, where the Hamiltonian describing the dynamics is quadratic in creation and annihilation operators. This particular fermion system was shown to be classically simulable by DiVincenzo and Terhal [19] as an extension of Valiant's findings.

With the use of recently developed techniques, the key objective of this project is to examine local operator spreading and scrambling in models where the dynamics are analytically and numerically tractable. This is not possible in the generic unitary evolution of many-body systems due to an exponential growth of the Hilbert space with system size. By simulating the dynamics of scrambling, a full picture of unitary evolution in simulatable models can be constructed, in hopes of gaining an understanding of dynamics of generic unitary evolution and whilst examining the differences between the two cases.

## II. BACKGROUND THEORY

### A. Quantum Information and Computing

Quantum Information is built upon the concept of quantum bits (qubits for short), represented as a linear combination of states in the standard basis,  $|\psi\rangle = \alpha|0\rangle + \beta|1\rangle$ , where  $\alpha, \beta$  are complex probability amplitudes and the vectors,  $|0\rangle, |1\rangle$  are the computational basis states that form the standard basis.

This is easily extended to multi-party or composite systems of  $n$  qubits. The total Hilbert space,  $\mathcal{H}$  of such a system is defined as the tensor product of  $n$  subsystem Hilbert Spaces,

$$\mathcal{H} = \otimes_n \mathcal{H}_n = \mathcal{H}_{n-1} \otimes \mathcal{H}_{n-2} \otimes \cdots \otimes \mathcal{H}_0 \quad (1)$$

Then we may write the computational basis states of this system, as strings of qubits, using the tensor product [20];

$$|x_1\rangle \otimes |x_2\rangle \otimes \cdots \otimes |x_n\rangle \equiv |x_1 x_2 \dots x_n\rangle.$$

The evolution and dynamics of such systems can be represented via quantum circuits, made up of quantum logic gates acting upon the qubits of a system. Analogous to a classical computer which is comprised of logic gates that act upon bit-strings of information. In contrast, quantum logic gates are linear operators acting on qubits, often represented in matrix form. This allows us to decompose a complicated unitary evolution into a sequence of linear transformations. A common practice is to create diagrams of such evolutions, with each quantum gate having their own symbol, analogous to circuit diagrams in classical computation, allowing the creation of complicated quantum circuitry that can be directly mapped to a sequence of linear operators acting on a one or more qubits. Some example gate symbols can be seen in Fig. 1.

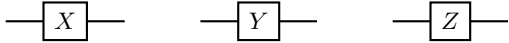


FIG. 1: Gate symbols for the Pauli operators in quantum circuits.

The gates shown in Fig. 1 are known as the Pauli operators, equivalent to the set of Pauli matrices,  $P \equiv \{X, Y, Z\}$  for which  $X, Y$  and  $Z$  are defined in their matrix representation as;

$$X = \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix}, \quad Y = \begin{bmatrix} 0 & i \\ -i & 0 \end{bmatrix}, \quad Z = \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} \quad (2)$$

These gates are all one-qubit gates, as they only act upon a single qubit. The simplest gate is that of the

Identity operator (or matrix),  $I$ , from which, an algebra can be formed with the Pauli matrices. Satisfying the following relations:

$$XY = iZ \quad YZ = iX \quad ZX = iY \quad (3)$$

$$YX = -iZ \quad ZY = -iX \quad XZ = -iY \quad (4)$$

$$X^2 = Y^2 = Z^2 = I \quad (5)$$

Notably, the set of Pauli matrices and the identity form the Pauli group,  $\mathcal{P}_n$ , defined as the  $4^n$   $n$ -qubit tensor products of the Pauli matrices (2) and the Identity matrix,  $I$ , with multiplicative factors,  $\pm 1$  and  $\pm i$  to ensure a legitimate group is formed. For clarity, consider the Pauli group on 1-qubit,  $\mathcal{P}_1$ ;

$$\mathcal{P}_1 \equiv \{\pm I, \pm iI, \pm X, \pm iX, \pm Y, \pm iY, \pm Z, \pm iZ\}, \quad (6)$$

From this, another group of interest can be defined, namely the Clifford group,  $\mathcal{C}_n$ , defined as a subset of unitary operators that normalise the Pauli group [21]. The elements of this group are the Hadamard, Controlled-Not and Phase operators.

The Hadamard,  $H$ , maps computational basis states to a superposition of computational basis states, written explicitly in it's action;

$$H|0\rangle = \frac{|0\rangle + |1\rangle}{\sqrt{2}},$$

$$H|1\rangle = \frac{|0\rangle - |1\rangle}{\sqrt{2}},$$

or in matrix form;

$$H = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 1 \\ 1 & -1 \end{bmatrix}$$

Controlled-Not ( $CNOT$ ) is a two-qubit gate. One qubit acts as a control for an operation to be performed on the other qubit. It's matrix representation is

$$CNOT = \begin{bmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{bmatrix} \quad (7)$$

The Phase operator, denoted  $S$  is defined as,

$$S = \begin{bmatrix} 1 & 0 \\ 0 & e^{i\frac{\pi}{2}} \end{bmatrix}$$

The  $CNOT$  operator is often used to generate entangled state. One such state is the maximally entangled 2-qubit state, called a Bell state,  $|\Phi^+\rangle = (|00\rangle +$

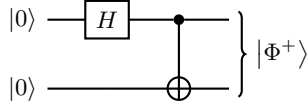


FIG. 2: Preparation of a Bell state from  $|0\rangle$  using a Hadamard and CNOT.

$|11\rangle)/\sqrt{2}$ . This is prepared from a  $|00\rangle$  state, by applying a Hadamard to the first qubit, and subsequently a Controlled-Not gate:

$$H \otimes I |00\rangle = \left( \frac{|0\rangle + |1\rangle}{\sqrt{2}} \right) |0\rangle \quad (8)$$

$$CNOT \left( \frac{|0\rangle + |1\rangle}{\sqrt{2}} \right) |0\rangle = \frac{|00\rangle + |11\rangle}{\sqrt{2}} \quad (9)$$

The corresponding circuit representation of this preparation is given in Fig. 2.

The output Bell state, is entangled and cannot be written in product form. For clarity, consider the following expression:

$$\begin{aligned} |\Phi^+\rangle &\neq [\alpha_0|0\rangle + \beta_0|1\rangle] \otimes [\alpha_1|0\rangle + \beta_1|1\rangle] \\ &\neq \alpha_0\beta_0|00\rangle + \alpha_0\beta_1|01\rangle + \alpha_1\beta_0|10\rangle + \alpha_1\beta_1|11\rangle \end{aligned}$$

the  $\alpha_0$  or  $\beta_1$  must be zero in order to ensure the  $|01\rangle$ ,  $|10\rangle$  vanish. However, this would make the coefficients of the  $|00\rangle$  or  $|11\rangle$  terms zero, breaking the equality. Thus,  $|\Phi^+\rangle$  cannot be written in product form and is said to be entangled. This defines a general condition for a arbitrary state to be entangled. [22]

### III. CLASSICALLY SIMULABLE QUANTUM CIRCUITS

#### A. Clifford Circuits and Stabilizer Formalism

Quantum circuits comprised of only quantum gates from the Clifford group are known as *Clifford Circuits*. This family of circuits are of considerable interest due to their classical simulability, that is, circuits of this construction can be efficiently simulated on a classical computer with polynomial effort via the Gottesman-Knill theorem. For this reason, Clifford gates do not form a set of universal quantum gates, meaning a universal quantum computer cannot be constructed using a set of Clifford gates. A set of universal quantum gates allows for any unitary operation to be approximated to arbitrary accuracy by a quantum circuit, constructed using only the original set of gates.

**Stabilizer Formalism.** An arbitrary, pure quantum state,  $|\psi\rangle$  is *stabilized* by a unitary operator,  $S$  if  $|\psi\rangle$  is an eigenvector of  $S$ , with eigenvalue 1, satisfying:

$$S|\psi\rangle = |\psi\rangle \quad (10)$$

It will be convenient to utilise the language of stabilizers, to define an initial state by the operators that stabilize it [23]. This can be seen from considering the Pauli matrices, and the unique states they stabilize. In the one-qubit case these are the +1 eigenstates of the Pauli matrices (omitting normalisation factors);

$$X(|0\rangle + |1\rangle) = |0\rangle + |1\rangle \quad (11)$$

$$Y(|0\rangle + i|1\rangle) = |0\rangle + i|1\rangle \quad (12)$$

$$Z|0\rangle = |0\rangle \quad (13)$$

If given a group or subgroup of unitaries,  $\mathcal{U}$ , the vector space,  $V$ , of  $n$  qubit states is stabilized by  $\mathcal{U}$  if every element of  $V$  is stable under action from every element of  $\mathcal{U}$ . This description is more appealing, as we can exploit mathematical techniques from group theory to describe quantum states and vector spaces. More precisely, the group  $\mathcal{U}$  can be described using its generators. In general, a set of elements  $g_1, \dots, g_d$  of a group,  $\mathcal{G}$ , generate the group if every element of  $\mathcal{G}$  can be written as a product of elements from the set of generators [24],  $g_1, \dots, g_d$ , such that  $\mathcal{G} := \langle g_1, \dots, g_d \rangle$ .

This allows for the description of a quantum state, and subsequently its dynamics, in terms of the generators of a stabilizer group. To see how the dynamics of a state are represented in terms of generators, consider a stabilizer state under the action of an arbitrary unitary operator:

$$UM|\psi\rangle = UMU^\dagger U|\psi\rangle \quad (14)$$

The state  $|\psi\rangle$  is an eigenvector of  $M$  if and only if,  $U|\psi\rangle$  is an eigenvector of  $UMU^\dagger$ . Thus, the application of an unitary operator transforms  $M \rightarrow UMU^\dagger$ . Moreover, if the state  $|\psi\rangle$  is stabilized by  $M$ , then the evolved state  $U|\psi\rangle$  will be stabilized by  $UMU^\dagger$ . If  $M$  is an element of a stabilizer group  $\mathcal{S}$  such that  $M_1, \dots, M_l$  generate  $\mathcal{S}$ , then  $UM_1U^\dagger, \dots, UM_lU^\dagger$  must generate  $USU^\dagger$ . This implies that to compute the dynamics of a stabilizer, only the transformation of the generators needs to be considered [25]. It is because of this, Clifford circuits are able to be efficiently classically simulated via the Gottesman-Knill theorem:

**Theorem 1** (Gottesman-Knill Theorem [26]). *Given an  $n$  qubit state  $|\psi\rangle$ , the following statements are equivalent:*

- $|\psi\rangle$  can be obtained from  $|0\rangle^{\otimes n}$  by CNOT, Hadamard and phase gates only.

- $|\psi\rangle$  is stabilized by exactly  $2^n$  Pauli operators
- $|\psi\rangle$  can be uniquely identified by the group of Pauli operators that stabilize  $|\psi\rangle$ .

Evolution via unitaries from the Clifford group, take an appealing form when acting on simple Pauli operators that act as stabilizers. In particular, Clifford unitaries map pauli operators to products of pauli operators, and since any product of Pauli operators can be written in terms of only products of  $X$  and  $Z$ , a generic Pauli string may be written as,

$$S \propto X_1^{v_{1x}} Z_1^{v_{1z}} \dots X_N^{v_{Nx}} Z_N^{v_{Nz}} \quad (15)$$

and can be specified by a binary vector,

$$\vec{v} = (v_{1x}, v_{1z}, \dots, v_{Nx}, v_{Nz}). \quad (16)$$

Thus, conjugation by Clifford unitaries on  $S$ , corresponds to local updates in  $\vec{v}$ . \*EXAMPLES\* The operation of Clifford unitaries on Pauli operators is shown in Table.

Unitary	Operation
CNOT	$X_1 \rightarrow X_1 X_2$
	$X_2 \rightarrow X_2$
	$Z_1 \rightarrow Z_1$
	$Z_2 \rightarrow Z_1 Z_2$
H	$X \rightarrow Z$
	$Z \rightarrow X$
S	$X \rightarrow X$
	$Z \rightarrow -Z$

## B. Non-interacting Fermi Circuits

### 1. Fermionic Fock Space

Another class of circuits that are classically simulable, are known as non-interacting or free fermi circuits. These circuits are built from a fermionic Fock space formalism, where the basic units of information are local fermionic modes (LFMs). Each mode can be occupied or unoccupied in the same fashion as qubits, where a mode  $j$  is associated with an occupation number,  $n_j = 0$  (unoccupied mode) or  $n_j = 1$  (occupied mode). The state space of a many-fermion system is the fermionic Fock space,  $\mathcal{F} = \mathcal{H}_0 + \mathcal{H}_1$ , where  $\mathcal{H}_0$  and  $\mathcal{H}_1$  correspond to subspaces of even and odd number of particles respectively. The natural basis vectors for a Fock space of  $m$  LFMs are Fock states  $|n_0, \dots, n_{m-1}\rangle$ , where each element is the previously mentioned occupation number,  $n_j$  at each site [27]. This system is similar to the qubit set-up as the Hilbert space for  $m$  LFMs is identified with the Hilbert space of  $m$  qubits such that

$$|n_1, \dots, n_{k-1}\rangle \equiv |n_0\rangle \otimes |n_1\rangle \otimes \dots \otimes |n_{m-1}\rangle, \quad n_j = 0, 1$$

where the RHS of the expression describes a system of  $m$  qubits. This allows a mapping of a qubit system (spin 1/2 system) onto a fermionic system via a Jordan-Wigner Transformation, such that the representation and dynamics of a quantum state are described using creation and annihilation operators from second quantization. The creation operator,  $a_j^\dagger$  creates a fermion at mode  $j$  (if the mode is unoccupied), whilst the annihilation operator,  $a_j$  removes a fermion at site  $j$  if the mode is occupied. In the case of fermions, these operators obey canonical anti-commutation relations:

$$\{a_i, a_j\} \equiv \{a_i^\dagger, a_j^\dagger\} = 0 \quad \{a_i, a_j^\dagger\} = \delta_{ij} I \quad (17)$$

The occupation number,  $n_j$  is then the eigenvalue of the number operator,  $\hat{n}_j = a_j^\dagger a_j$ . The creation and annihilation operators act on Fock states in the following way:

$$\begin{aligned} a_j^\dagger |n_0, \dots, n_j, \dots, n_{m-1}\rangle &= (-1)^{\sum_{s=0}^{j-1} n_s} \delta_{n_j, 1} |n_0, \dots, n'_j, \dots, n_{m-1}\rangle \\ a_j |n_0, \dots, n_j, \dots, n_{m-1}\rangle &= (-1)^{\sum_{s=0}^{j-1} n_s} \delta_{n_j, 0} |n_0, \dots, n'_j, \dots, n_{m-1}\rangle \end{aligned}$$

Where  $n'_j$  is the updated occupation number at the  $j$ th site. To prepare an arbitrary Fock state,  $|\mathbf{n}\rangle$ , from a

vacuum,  $|\mathbf{0}\rangle = |0, 0, \dots, 0\rangle$ , with fermions occupying arbitrary positions, creation operators are labeled to act

on specified modes:

$$|\mathbf{n}\rangle = a_{i_1}^\dagger a_{i_2}^\dagger \dots a_{i_m}^\dagger |\mathbf{0}\rangle \quad (18)$$

## 2. Fermionic Circuits

A fermionic circuit is constructed from a sequence of elementary gates that mediate an interaction between modes  $j$  and  $k$ , written as  $U = e^{iH_g} H_g$  is a general gate Hamiltonian:

$$H_g = b_{jj} a_j^\dagger a_j + b_{kk} a_k^\dagger a_k + b_{jk}^* a_k^\dagger a_j, \quad (19)$$

where  $b_{jj}, b_{jk}, b_{kk}$  are complex coefficients that form a hermitian matrix,  $\mathbf{b}$ . The circuit is constructed by defining the sequence of gates as  $U = U_{\text{poly}(n)} \dots U_2 U_1$ , such that the number of gates is no more than a polynomial in order to be simulated classically. Then considering the evolution of Fock state in a circuit, given that the circuit preserves the number of fermions ( $U|\mathbf{0}\rangle = |\mathbf{0}\rangle$ ), the state evolves as:

$$U a_j^\dagger |\mathbf{0}\rangle = U a_j^\dagger U^\dagger U |\mathbf{0}\rangle = U a_j^\dagger U^\dagger |\mathbf{0}\rangle \quad (20)$$

Where  $U$  acts by conjugation as

$$U a_j^\dagger U^\dagger = \sum_s B_{is} a_s^\dagger \quad (21)$$

and the matrix  $B = \exp(i\mathbf{b})$  is defined from the previous coefficient matrix in Eq. (19). To compute the final state of the system, only the matrix  $B$  needs evaluation, which is classically simulable if  $U$  contains a polynomial number of gates [19]. More generally, if  $U$  acts on an arbitrary Fock state,  $|\mathbf{n}\rangle$ , it's action is written as:

$$U|\mathbf{n}\rangle = U a_{i_1}^\dagger a_{i_2}^\dagger \dots a_{i_m}^\dagger |\mathbf{0}\rangle \quad (22)$$

$$= \sum_{j_1 \dots j_k} V_{i_1, j_1} \dots V_{i_k, j_k} a_{j_1}^\dagger a_{j_2}^\dagger \dots a_{j_k}^\dagger |\mathbf{0}\rangle \quad (23)$$

Some simpler operators that act on Fock states, are the spin 1/2 Pauli operators, that are mapped onto a non-interacting fermion system via the Jordan Wigner Transformation (JWT). In particular, the set of Pauli operators that act on the  $j$ th mode,  $\{X_j, Y_j, Z_j\}$ , can be expressed in terms of the creation and annihilation operators by first defining the operators  $\sigma_j^\pm = \frac{1}{2}(X_j \pm iY_j)$ . Then in terms of the creation and annihilation operators:

$$\sigma_j^+ = e^{i\pi \sum_{m=0}^{j-1} a_m^\dagger a_m} a_j^\dagger \quad (24)$$

$$\sigma_j^- = e^{i\pi \sum_{m=0}^{j-1} a_m^\dagger a_m} a_j \quad (25)$$

A transform of the creation and annihilation operators can also be made, to define the Majorana fermion operators that effectively split each LFM into halves [28];

$$c_{2j} = a_j + a_j^\dagger \quad c_{2j+1} = \frac{a_j - a_j^\dagger}{i} \quad (26)$$

Majorana fermion operators satisfy the anti-commutation relation;

$$\{c_j, c_k\} = 2\delta_{jk} \quad (27)$$

The majorana fermion operators may also be expressed in terms of Pauli operators,

$$c_{2j} = X_j \prod_{k=0}^{j-1} Z_k \quad c_{2j+1} = Y_j \prod_{k=0}^{j-1} Z_k \quad (28)$$

where the term  $\prod_{k=0}^{j-1} Z_k$  is called a Jordan Wigner transformation string.

## IV. OPERATOR SCRAMBLING

### A. Operator Spreading

As previously mentioned, quantum scrambling is the process in which local information encoded by a simple product operator is rapidly spread over a large number of degrees of freedom via a unitary time evolution. The information becomes highly non-local and the initial simple product operator becomes a complicated sum of product operators. To give picture of operator spreading consider a local operator,  $O_j$ , which has support on either one site or a finite region of sites, so that it acts like the Identity on sites without support. The simplest local operator would be a Pauli operator acting on a single site,  $j$ , e.g  $O_j \equiv X_j$ . Then the evolution of  $O_j$ , is given, in the heisenberg picture, by  $O_j(t) = U^\dagger(t) O_j U(t)$ , where  $U(t)$  is some unitary time evolution operator. Following this evolution, the local operator with minimal support,  $O_j$  has evolved to  $O_j(t)$  with support over a large region of sites [10]. Operators that grow in this way, will spread ballistically [29–31] and are often characterised by the out-of-time ordered correlator (OTOC) [32] and the square-commutator[14]. This also gives an intuitive picture of operator spreading, with the squared commutator defined as

$$C(t) = \langle [O(t), W_i][O(t), W_i]^\dagger \rangle \quad (29)$$

where  $W_i$  is a static local operator at site  $i$ [33]. Then at  $t = 0$ ,  $O(t)$  acts on a single site or a finite region of

sites, such that it commutes with the static operator,  $V_i$  and  $C(t) = 0$ . Once the operator spreads, and becomes more non-local, the commutator increases as it's support overlaps with  $V_i$ .

## B. Entanglement Growth in Operator Space

Under a generic unitary evolution, we expect an initially simple product operator, for example, a Pauli string,  $O = P_1, \dots, P_n$  to evolve into a linear superposition of Pauli strings [34, 35]. This entanglement growth corresponds to an exponential increase in the possible number of Pauli strings, requiring the use of stabilizer states from Clifford circuits. However, since Clifford unitaries only map products of Pauli operators to products of Pauli operators, a Pauli string will remain unentangled [36]. Despite this, the scrambling behaviour can be recovered, by the construction of *super-Clifford operators*, whilst remaining classically simulable as shown by Blake and Linden [37].

### 1. Blake and Linden's Construction

Blake and Linden introduce a family of circuits that exhibit scrambling on the space spanned by Pauli operators. They present a gate-set of 'super-clifford operators' that generate a near-maximal amount of operator entanglement within the Pauli operator space, called super-clifford circuits, when under time-evolution. These super-clifford operators remain classically simulable, via an extension of stabilisers to operator space, called 'super-stabilisers'. Super-Clifford operators 'act' on Pauli operators, via conjugation in the Heisenberg picture. Namely, the first super-Clifford operator in the Gate set is denoted as **Z.H**, which is identical to conjugation by a Phase gate,  $T$ ;

$$T^\dagger X T = \frac{X - Y}{\sqrt{2}} \quad T^\dagger Y T = \frac{X + Y}{\sqrt{2}} \quad (30)$$

Following this, the operators,  $X$  and  $Y$  can be changed to a state-like representation, with  $X$  denoted as  $[0]$  and  $Y$  denoted as  $[1]$ . The second gate in the set of super-Clifford operators, is the **SWAP** gate,

$$\mathbf{SWAP}[01] = [10] \quad (31)$$

formed from the regular 2-qubit SWAP gate, that swaps two nearest neighbour qubits. It conjugates Pauli operators in the following way:

$$\mathbf{SWAP}^\dagger X_1 Y_2 \mathbf{SWAP} = Y_1 X_2 \quad (32)$$

The third gate, denoted **C3**, acts as a combination of controlled-**Y** super-operators,

$$\mathbf{C3}[000] = \mathbf{CY}_{12}\mathbf{CY}_{13}[000] = [000] \quad (33)$$

$$\mathbf{C3}[100] = \mathbf{CY}_{12}\mathbf{CY}_{13}[100] = -[111] \quad (34)$$

These three super-operators form the gate set  $\{\mathbf{SWAP}, \mathbf{Z.H}, \mathbf{C3}\}$ , which generates entanglement through unitary evolution in operator space, as this gate set maps Pauli strings to a linear superposition of Pauli strings. Even though the a simple string of Pauli operators can evolve into a sum of potentially exponential operator strings, the dynamics can be computed classically by extending the formalism of stabilizer states to operator space. Following this, Blake and Linden showed that the gate set was capable of generating near-maximal amounts of entanglement (slightly less than the Page value [38]) among Pauli strings on a chain of 120 qubits, quantified by the von Neumann entropy. This was only shown for operators with global support, and therefore shows no notion of operator spreading or entanglement growth in local operators.

### 2. Scrambling in Free-Fermi Circuits

As previously mentioned, a spin 1/2 or qubit, system can be mapped onto a fermionic system via a Jordan Wigner transformation, with the introduction of Majorana fermion operators. The evolution of Majorana fermion Operators has been shown to exhibit operator entanglement growth by Prosen and Pižorn [39] by mapping an Ising spin-1/2 system onto a Majorana fermion representation. Further efforts from now on will be directed towards simulating non-interacting fermion systems and computing the entanglement entropy so that the dynamics of local operator spreading and entanglement growth can be fully investigated.

## V. FURTHER DIRECTION

Further research will be directed towards free-fermion systems and how to simulate them. This will be carried out by writing a python program to simulate random unitary circuits and computing the entanglement entropy. We will first reproduce Blake and Linden's results for the  $N = 120$  qubits case of super-Clifford evolution, and then go on to simulate free-fermion circuits to confirm a non-trivial operator scrambling. Following this, the idea of combining Clifford and non-interacting fermion circuits will be explored to investigate if any amplification of the scrambling occurs [40], along with

relaxing the condition of nearest-neighbour interactions,  
by using long-ranged gates.

- 
- [1] A. Polkovnikov, K. Sengupta, A. Silva, and M. Vengalattore, “Nonequilibrium dynamics of closed interacting quantum systems,” *Reviews of Modern Physics*, vol. 83, pp. 863–883, aug 2011.
  - [2] P. Calabrese and J. Cardy, “Entanglement entropy and conformal field theory,” *Journal of Physics A: Mathematical and Theoretical*, vol. 42, p. 504005, dec 2009.
  - [3] D. Harlow, “Jerusalem lectures on black holes and quantum information,” *Rev. Mod. Phys.*, vol. 88, p. 015002, Feb 2016.
  - [4] K. Jensen, “Chaos in  $\text{ads}_2$  holography,” *Physical Review Letters*, vol. 117, sep 2016.
  - [5] L. Susskind, “Why do things fall?,” 2018.
  - [6] Y. Sekino and L. Susskind, “Fast scramblers,” *Journal of High Energy Physics*, vol. 2008, pp. 065–065, oct 2008.
  - [7] S. H. Shenker and D. Stanford, “Black holes and the butterfly effect,” *Journal of High Energy Physics*, 2014.
  - [8] M. Nozaki, T. Numasawa, and T. Takayanagi, “Quantum entanglement of local operators in conformal field theories,” *Physical Review Letters*, vol. 112, mar 2014.
  - [9] J. Maldacena, S. H. Shenker, and D. Stanford, “A bound on chaos,” *Journal of High Energy Physics*, vol. 2016, aug 2016.
  - [10] V. Khemani, A. Vishwanath, and D. A. Huse, “Operator spreading and the emergence of dissipative hydrodynamics under unitary evolution with conservation laws,” *Physical Review X*, vol. 8, sep 2018.
  - [11] C. W. von Keyserlingk, T. Rakovszky, F. Pollmann, and S. L. Sondhi, “Operator hydrodynamics, otocs, and entanglement growth in systems without conservation laws,” *Phys. Rev. X*, vol. 8, p. 021013, Apr 2018.
  - [12] T. Rakovszky, F. Pollmann, and C. W. von Keyserlingk, “Diffusive hydrodynamics of out-of-time-ordered correlators with charge conservation,” *Phys. Rev. X*, vol. 8, p. 031058, Sep 2018.
  - [13] S. Grozdanov, K. Schalm, and V. Scopelliti, “Black hole scrambling from hydrodynamics,” *Physical Review Letters*, vol. 120, jun 2018.
  - [14] M. Blake, H. Lee, and H. Liu, “A quantum hydrodynamical description for scrambling and many-body chaos,” *Journal of High Energy Physics*, vol. 2018, oct 2018.
  - [15] Y. Li, X. Chen, and M. P. A. Fisher, “Quantum zeno effect and the many-body entanglement transition,” *Phys. Rev. B*, vol. 98, p. 205136, Nov 2018.
  - [16] Y. Li, R. Vasseur, M. P. A. Fisher, and A. W. W. Ludwig, “Statistical mechanics model for clifford random tensor networks and monitored quantum circuits,” 2021.
  - [17] T. Farshi, J. Richter, D. Toniolo, A. Pal, and L. Masanes, “Absence of localization in two-dimensional clifford circuits,” 2022.
  - [18] L. G. Valiant, “Quantum computers that can be simulated classically in polynomial time,” in *STOC ’01*, 2001.
  - [19] B. M. Terhal and D. P. Divincenzo, “Classical simulation of noninteracting-fermion quantum circuits Typeset using REVTeX 1,” tech. rep., 2001.
  - [20] B. Schumacher and M. Westmoreland, *Quantum Processes Systems, and Information*. Cambridge University Press, 2010.
  - [21] A. R. Calderbank, E. M. Rains, P. W. Shor, and N. J. A. Sloane, “Quantum error correction via codes over  $\text{gf}(4)$ ,” 1996.
  - [22] M. A. Nielsen and I. L. Chuang, *Quantum Computation and Quantum Information: 10th Anniversary Edition*. Cambridge University Press, 2010.
  - [23] S. Aaronson and D. Gottesman, “Improved simulation of stabilizer circuits,” *Phys. Rev. A*, vol. 70, p. 052328, Nov 2004.
  - [24] A. R. Calderbank, E. M. Rains, P. W. Shor, and N. J. A. Sloane, “Quantum error correction and orthogonal geometry,” *Physical Review Letters*, vol. 78, pp. 405–408, jan 1997.
  - [25] D. Gottesman, “Theory of fault-tolerant quantum computation,” *Physical Review A*, vol. 57, pp. 127–137, jan 1998.
  - [26] D. Gottesman, “The heisenberg representation of quantum computers,” 1998.
  - [27] G. De Felice, A. Hadzihanovic, and K. F. Ng, “A diagrammatic calculus of fermionic quantum circuits,” *Logical Methods in Computer Science*, vol. 15, no. 3, 2019.
  - [28] S. B. Bravyi, L. D. Landau, and A. Y. Kitaev, “Fermionic quantum computation,” tech. rep., 2000.
  - [29] D. A. Roberts, D. Stanford, and L. Susskind, “Localized shocks,” *Journal of High Energy Physics*, vol. 2015, mar 2015.
  - [30] E. H. Lieb and D. W. Robinson, “The finite group velocity of quantum spin systems,” *Commun. Math. Phys.*, vol. 28, pp. 251–257, 1972.
  - [31] T. Schuster and N. Y. Yao, “Operator growth in open quantum systems,” 2022.
  - [32] S. Xu and B. Swingle, “Scrambling Dynamics and Out-of-Time Ordered Correlators in Quantum Many-Body Systems: a Tutorial,” feb 2022.
  - [33] X. Chen and T. Zhou, “Operator scrambling and quantum chaos,” 2018.
  - [34] D. A. Roberts, D. Stanford, and A. Streicher, “Operator growth in the SYK model,” *Journal of High Energy Physics*, vol. 2018, jun 2018.
  - [35] A. Nahum, J. Ruhman, S. Vijay, and J. Haah, “Quantum entanglement growth under random unitary dynamics,” *Physical Review X*, vol. 7, jul 2017.
  - [36] A. Nahum, S. Vijay, and J. Haah, “Operator spreading in random unitary circuits,” *Physical Review X*, vol. 8, apr 2018.
  - [37] M. Blake and N. Linden, “Quantum circuits with classically simulable operator scrambling,” feb 2020.
  - [38] D. N. Page, “Average entropy of a subsystem,” *Physical Review Letters*, vol. 71, pp. 1291–1294, aug 1993.
  - [39] T. Prosen and I. Pizorn, “Operator space entanglement entropy in a transverse ising chain,” *Physical Review A*, vol. 76, sep 2007.



- 
- [40] R. Jozsa and A. Miyake, “Matchgates and classical simulation of quantum circuits,” apr 2008.