

Thermodynamic fragments for magic states in quantum computation

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[To be sharpened] Magic states are key ingredients in schemes to realise universal fault-tolerant quantum computation. Theories of magic states attempt to quantify this computational element via monotones and determine how these states may be efficiently transformed into useful forms. Here we introduce the concept of ‘fragments’, which generalise the concept of magic monotones and has a natural thermodynamic structure based on majorisation. From this perspective magic can be viewed as a form of free energy within each fragment and is constrained by relative majorisation relations but now on quasi-probability distributions. Notably this approach allows us to incorporate actual physical constraints, for example noise models with particular bias or temperature-dependent features, and study how these constrain general magic distillation protocols. In this context we present general temperature-dependent bounds on distillation rates that any theory of magic must respect. Significantly, this analysis also presents a thermodynamic context which cannot be analysed via traditional methods based on thermodynamic entropies, due to the presence of negativity, and raises novel questions in the context of statistical mechanics.

I. INTRODUCTION AND BACKGROUND

1. Fault-tolerance [1–9]
2. Magic [10–17]
3. Bringing in majorisation [18–26]

A. Introduction

Magic states are necessary for achieving universal quantum computation within fault-tolerant schemes. Identifying magic as a resource for quantum universality has led to several theories which try to provide a framework for its quantification and manipulation [CITE]. The main question that such a theory attempts to answer is:

Given two magic states ρ and ρ' is there a free operation that can convert ρ to ρ' ?

We are interested in all resource theories of magic $\mathcal{R} = (\mathcal{F}, \mathcal{O})$ in which free operations cannot generate any amount of resource. Further denote by \mathcal{D} the set of states considered under the theory, that is the union of free and resource states. The structure of such theory is described by a partial order [CITE], hereinafter called a *pre-order*, $\prec_{\mathcal{R}}$ between states. We write $\rho' \prec_{\mathcal{R}} \rho$ iff there exists $\mathcal{E} \in \mathcal{O}$ such that $\mathcal{E}(\rho) = \rho'$. Naturally, states may be incomparable under the given theory, meaning that there exists no free operation that converts one to the other. We further call $\mathcal{R}' = (\mathcal{F}', \mathcal{O}')$ a *subtheory* of \mathcal{R} iff

$\mathcal{F}' \subseteq \mathcal{F}$ and $\mathcal{O}' \subseteq \mathcal{O}$. The above notation will be used for general resource theories as well.

Formally, the no resource generation condition on the theories translates into two assumptions:

I Free operations send free states into free states, $\mathcal{E} : \mathcal{F} \mapsto \mathcal{F}$, for all $\mathcal{E} \in \mathcal{O}$;

II Resource theory \mathcal{R} is a completely free state preserving theory, in the sense that for any d -dimensional ancilla system and all free operations \mathcal{E} , $(\mathbb{1}_d \otimes \mathcal{E})\sigma \in \mathcal{F}$ whenever $\sigma \in \mathcal{F}$.

The first assumption simply states that resources cannot be generated for free and is a minimal requirement for a resource theory. An immediate consequence is that if statistical mixing is included in \mathcal{O} , then the resource theory is convex. Convex resource theories have attracted a lot of attention recently [CITE] and include the magic theories discussed in Section IB. The second assumption implies that resources cannot be generated even when ancillas are allowed [example of T state generation on Bell state by Campbell].

Monotones are often used [CITE] to address the question of state convertibility, although such approaches are usually generic.

The monotonicity condition reflects the no resource generating property of free operations, so that monotones respect the pre-order $\prec_{\mathcal{R}}$ of the theory. A monotone of any general resource theory is a projection of the theory onto the non-negative real numbers, collapsing the pre-order of the theory to the total order defined on the real line. Our contribution is the introduction of a generalised notion of *resource projection* which maps a general resource theory onto a subtheory which in principle still retains a partial structure (as opposed to the real line). Applying this notion on existing magic theories highlights

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the hidden stochasticity that governs magic state conversions. We show that a magic theory can be subdivided into *fragments* [\[expand\]](#)

B. Previous work

[\[Section needs a lot of fixing\]](#)

The stabiliser theory of magic comprises of the so-called “stabiliser” states (STAB) and operations (SO), while non-stabiliser (resource) states are called magic. The stabiliser operations can be expressed in terms of a Stinespring dilation as

$$\mathcal{E}(\rho) = \text{tr}_E[U(\rho \otimes \sigma_E)U^\dagger], \quad (1)$$

for an ancilla stabiliser state σ_E . The motivation of the theory stems from the fact that stabiliser operations are experimentally straightforward to implement and they can be used to detect and correct errors on the stabiliser states due to their construction [\[CITE\]](#). The Gottesman-Knill theorem however indicates that stabiliser operations need to be supplemented with magic states in order to achieve universality, justifying the term “magic”.

Generalisations of the stabiliser theory appear in the literature intending to include broader classes of operations [\[CITE\]](#). The class of stabiliser preserving operations (SPO) is defined as the set of CPTP maps that send stabiliser states into stabiliser states [\[27\]](#). An important subclass of SPO is the set of completely stabiliser preserving operations (CSPO), which intuitively cannot induce “non-stabiliserness” even when applied to only part of a quantum state, i.e. operations \mathcal{E} such that $(\mathbb{1}_d \otimes \mathcal{E})\sigma \in \text{STAB}$ for all positive dimensions d whenever $\sigma \in \text{STAB}$.

Even though non-stabiliserness is a necessary resource for universality, it has been proven insufficient for magic state distillation [\[15, 28\]](#). In fact, all states with non-negative Wigner distributions have been proven to be efficiently classically simulable in [\[29\]](#), a result that serves as a generalization of the Gottesman-Knill theorem. The Wigner distribution of a state in odd prime dimensions is discussed rigorously in Section [IIB](#) and arises as the unique quasi-probability representation of quantum theory that identifies non-contextuality exactly with the states that are efficiently classically simulable [\[11, 30\]](#). In this framework, the stabiliser states are the only pure states represented with non-negative distributions [\[16\]](#). However, there exist mixed states with non-negative Wigner distributions that are not mixtures of stabiliser states [\[4\]](#). Therefore, stabiliser-preserving theories have been extended to a theory that preserves state “Wigner positivity” [\[12\]](#), formally defined in Section [IIB](#) for odd prime dimensions. Informally, it can be considered as the maximal theory of magic $\mathcal{R}_{\text{max}} = (\mathcal{F}_{\text{max}}, \mathcal{O}_{\text{max}})$, where free states have non-negative Wigner distributions and free operations completely preserve this property.

[Things we MUST emphasize:](#)

1. Perhaps a nice lead-in question: “What happens if we view stabiliser states as thermodynamic equilibrium states and magic as a form of free energy?”
2. We have found a scenario in which it is impossible to describe a thermodynamic structure using any entropic approach!
3. We can tackle more ‘physicsy’ questions like: how much magic can be distilled via available operations with some given fixed-point structure?
4. This allows a diagnostic on the kind of operations needed to do good distillation. I.e. what fixed point structure should they have?
5. We go beyond the concept of monotones and replace a monotone with a σ -fragment.
6. We can get both upper and lower bounds on magic distillation.

II. PRELIMINARIES AND DEFINITIONS

A. Stabilizer Theory

Let $\{|k\rangle\}_{k \in \mathbb{Z}_d}$ be the standard computational basis, defined over $\mathbb{Z}_d = \{0, 1, \dots, d-1\}$, with d an odd prime. Since the field has character d , addition and multiplication on the field are always considered modulo d . The Hilbert space of any system associated with this scheme is $\mathcal{H}_d := \text{span}\{|k\rangle : k \in \mathbb{Z}_d\}$.

The generalised Pauli matrices X, Z can be defined by their respective roles as shift and phase operators,

$$X|k\rangle = |k+1\rangle \quad (2)$$

$$Z|k\rangle = \omega^k |k\rangle, \quad (3)$$

where $\omega := e^{2\pi i/d}$ is the d -th root of unity and addition is modulo d . The Hilbert space \mathcal{H}_d is associated with a phase space $\mathcal{P}_d = \mathbb{Z}_d \times \mathbb{Z}_d$, where every point $\mathbf{x} := (x, p)$ corresponds to a displacement operator, defined as

$$D_{\mathbf{x}} := \tau^{xp} X^x Z^p, \quad (x, p) \in \mathcal{P}_d. \quad (4)$$

The phase factor $\tau := -\omega^{1/2}$ ensures unitarity. For a composite system with product dimension $d = d_1 \dots d_n$ and Hilbert space, $\mathcal{H}_d = \mathcal{H}_{d_1} \otimes \dots \otimes \mathcal{H}_{d_n}$, the displacement operators are defined as

$$D_{\mathbf{x}} := D_{(x_1, p_1)} \otimes \dots \otimes D_{(x_n, p_n)}, \quad (5)$$

where

$$\mathbf{x} := (x_1, p_1, x_2, p_2, \dots, x_n, p_n) \in \mathcal{P}_{d_1} \times \dots \times \mathcal{P}_{d_n} =: \mathcal{P}_d,$$

denotes the phase space point for an n -copy system.

The displacement operators form the Heisenberg-Weyl group [CITE] under matrix multiplication modulo phases,

$$\text{HW}_d^n := \{\tau^k D_{\mathbf{x}} : k \in \mathbb{Z}_d, \mathbf{x} \in \mathcal{P}_d^n\}. \quad (6)$$

The Clifford operations \mathcal{C}_d^n are then defined as the set of unitaries that normalise the Heisenberg-Weyl group. We may define the pure stabiliser states as those states obtained by acting on $|0\rangle$ with Clifford unitaries. Finally, we define STAB as the convex hull of all pure stabiliser states, namely all probabilistic mixtures of states of the form $U|0\rangle\langle 0|U^\dagger$ where U is some Clifford unitary. The stabiliser theory forms a magic resource theory $\mathcal{R} = (\text{STAB}, \text{SO})$ and the free operations are stabiliser operations (SO), defined as compositions of preparation of computational basis states, Clifford unitaries, measurement in the computational basis, and the ability to discard subsystems.

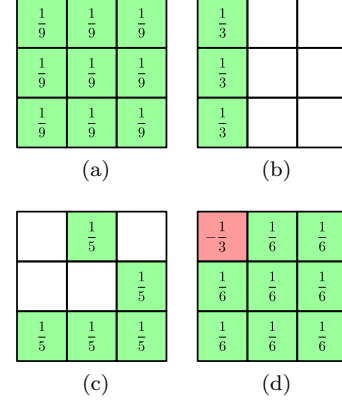


FIG. 1. **Qutrit Wigner distributions of varying magic.** (a) Maximally mixed state $\frac{1}{3}\mathbb{1}$; (b) Stabilizer zero state $|0\rangle\langle 0|$; (c) A non-stabiliser Wigner-positive state; (d) Magic Strange state $|S\rangle = \frac{1}{\sqrt{2}}(|1\rangle - |2\rangle)$, coined in [11].

B. Wigner representations for quantum states and quantum operations

We define the phase-point operators at every phase space point $\mathbf{x} \in \mathcal{P}_d$,

$$A_{\mathbf{x}} := \frac{1}{d} \sum_{\mathbf{z} \in \mathcal{P}_d} \omega^{\eta(\mathbf{x}, \mathbf{z})} D_{\mathbf{z}}, \quad (7)$$

where $\eta(\mathbf{x}, \mathbf{z})$ is the standard symplectic form, given explicitly by

$$\eta(\mathbf{x}, \mathbf{z}) := \mathbf{z}^T \begin{pmatrix} 0 & \mathbb{1} \\ -\mathbb{1} & 0 \end{pmatrix} \mathbf{x}, \quad (8)$$

where $0, \mathbb{1}$ denote the $n \times n$ zero and identity matrices.

The phase-point operators form an orthogonal Hermitian operator basis as shown in Appendix A. Therefore, any quantum state $\rho \in \mathcal{B}(\mathcal{H}_d)$ can be expressed as a linear combination of them,

$$\rho = \sum_{\mathbf{z} \in \mathcal{P}_d} W_{\rho}(\mathbf{z}) A_{\mathbf{z}}, \quad (9)$$

where the coefficient vector W_{ρ} is the Wigner distribution of state ρ ,

$$W_{\rho}(\mathbf{x}) := \frac{1}{d} \text{tr}[A_{\mathbf{x}} \rho]. \quad (10)$$

For any quantum state ρ , the Wigner distribution $W_{\rho}(\mathbf{x})$ is readily seen to be a d^2 -dimensional quasi-probability distribution over \mathcal{P}_d (see Appendix A for details). In Fig. (1), we show Wigner distributions of different types of qutrit states.

We also have Wigner representations for general quantum channels. We may exploit the channel-state duality and use the normalised Choi-Jamiołkowski state

$$\mathcal{J}(\mathcal{E}) := \frac{1}{d_A} (\mathbb{1} \otimes \mathcal{E}) \sum_{i,j} |ii\rangle \langle jj| \quad (11)$$

to extend the definition of the Wigner distribution to a quantum channel $\mathcal{E} : \mathcal{B}(\mathcal{H}_A) \rightarrow \mathcal{B}(\mathcal{H}_B)$, via the expression

$$W_{\mathcal{E}}(\mathbf{y}|\mathbf{x}) := d_A^2 W_{\mathcal{J}(\mathcal{E})}(\bar{\mathbf{x}} \oplus \mathbf{y}) \quad (12)$$

$$= \frac{1}{d_B} \text{tr}_B[A_{\mathbf{y}} \mathcal{E}(A_{\mathbf{x}})], \quad (13)$$

where $\bar{\mathbf{x}} := (x, -p)$.

The specific form of Eq. (12) is chosen so that Wigner distributions of operations act as transition matrices for Wigner distributions of states, $W_{\mathcal{E}(\rho)} = W_{\mathcal{E}} W_{\rho}$. In particular, CPTP operations that map between density operators of equal dimensions and have non-negative Wigner distributions correspond to stochastic matrices, as shown in Appendix A

C. Magic monotones

We highlight the defining property of a monotone, as it is useful in later discussion.

Definition 1 (Resource monotone). *Let $\mathcal{R} = (\mathcal{F}, \mathcal{O})$ be a resource theory. A resource monotone \mathcal{M} is a projection from the set of quantum states of the theory onto the real line, so that \mathcal{M} is monotonically decreasing under free operations,*

$$\mathcal{M}(\rho_1) \leq \mathcal{M}(\rho_2) \text{ whenever } \rho_1 \prec_{\mathcal{R}} \rho_2. \quad (14)$$

[We can skip mana and just discuss sum-negativity]

One of the most fundamental and commonly used magic monotones is the *mana* of a state [11], defined as

$$\text{mana}(\rho) := \ln(2 \text{sn}(\rho) + 1), \quad (15)$$

where the *sum-negativity* (sn) [11] is the sum of the negative components in W_{ρ} ,

$$\text{sn}(\rho) := \sum_{\mathbf{x}: W_{\rho}(\mathbf{x}) < 0} |W_{\rho}(\mathbf{x})|. \quad (16)$$

Mana is an additive¹ magic monotone, so it provides an analytical, necessary condition for many-copy magic state interconvertibility.

D. Majorisation

Majorisation is a collection of powerful tools that has recently found many applications in quantum information theory [CITE]. It describes the [disorder / non-uniformity] of distributions that undergo stochastic transformations.

To formally state majorisation results, we first denote by $S_d(\mathbf{d})$ the set of $(d \times d)$ stochastic matrices that preserve the probability vector \mathbf{d} . Specifically, for any $S \in S_d(\mathbf{d})$, all matrix elements are non-negative, all rows sum to 1 and $S\mathbf{d} = \mathbf{d}$. The set $S_d(\mathbf{d})$ forms a group under matrix multiplication for all \mathbf{d} with positive components.

Majorisation describes quantum thermodynamics exactly in the absence of quantum coherence. The use of majorisation in this setting provides vital intuition for our purposes. At any given temperature β^{-1} , the thermal state γ_{β} is thermodynamically the most disordered state. Thermal operations are defined as operations that cannot extract energy from the Gibbs state, $\mathcal{E}(\gamma_{\beta}) = \gamma_{\beta}$. Convertibility between states via thermal operations is equivalent to a stochasticity condition on the energy level populations of the states [CITE]. Informally, there exists a thermal operation \mathcal{E} such that $\tau = \mathcal{E}(\rho)$ if and only if there exists a matrix $S \in S_d(\mathbf{d})$ such that $\mathbf{q} = S\mathbf{p}$, where \mathbf{q}, \mathbf{p} and \mathbf{d} are the energy level population vectors of $\tau, \rho, \gamma_{\beta}$ respectively. Formal statements of the above can be found [in the appendix].

Based on this setting, we define majorisation as follows.

Definition 2 (*d*-majorisation). *Given $\mathbf{x}, \mathbf{y}, \mathbf{d} \in \mathbb{R}^d$, such that the components of \mathbf{d} are positive, \mathbf{y} is said to *d*-majorise \mathbf{x} , iff there exists a matrix $S \in S_d(\mathbf{d})$ such that $\mathbf{x} = S\mathbf{y}$.*

We denote this pre-order by $\mathbf{x} \prec_d \mathbf{y}$. If $\mathbf{d} = \frac{1}{d}\mathbf{1}$, the *d*-dimensional uniform distribution, then $S_d(\mathbf{d})$ is the set of doubly stochastic matrices and we retrieve the familiar notion of majorisation from entanglement theory. [CITE]

The pre-order imposed by *d*-majorisation admits a numerically efficient reformulation in terms of Lorenz curves. Let the vector \mathbf{u}^{\downarrow} denote the vector $\mathbf{u} \in \mathbb{R}^d$ with its components arranged in non-increasing order.

Definition 3 (Lorenz curve). *Let $\mathbf{w}, \mathbf{d} \in \mathbb{R}^d$, where the components of \mathbf{d} are positive with $D = \sum_{i=1}^d d_i$ and denote by $\tilde{\mathbf{w}} := (w_1/d_1, \dots, w_d/d_d)^T$ the vector of ratios between the corresponding components of \mathbf{w} and \mathbf{d} .*

Finally, denote by $\pi : \mathbb{Z}_d \mapsto \mathbb{Z}_d$ the permutation that sorts $\tilde{\mathbf{w}}$, $(\tilde{\mathbf{w}}^{\downarrow})_i = w_{\pi(i)}$ for all $i = 1, \dots, d$.

Consider the piecewise linear curve obtained by joining the points $\{(0, 0)\} \cup \{(x_k, L_{\mathbf{w}|\mathbf{d}}(k))\}_{k=1, \dots, d}$, where

$$(x_k, L_{\mathbf{w}|\mathbf{d}}(k)) := \left(\frac{1}{D} \sum_{i=1}^k d_{\pi(i)}, \sum_{i=1}^k w_{\pi(i)} \right). \quad (17)$$

We define the set of points on this curve, $L_{\mathbf{w}|\mathbf{d}}(x)$, $x \in [0, 1]$, as the Lorenz curve of vector \mathbf{w} with respect to \mathbf{d} .

Components x_k are rescaled by D so that comparison of curves with unequal dimensions is possible. In fact, the Lorenz curves $L_{\mathbf{w}|\mathbf{d}}$ and $L_{\mathbf{w} \otimes \mathbf{d} | \mathbf{d} \otimes \mathbf{d}}$, where \otimes denotes the Kronecker product, are identical. Furthermore, a Lorenz curve $L_{\mathbf{w}|\mathbf{d}}(x)$ is always concave in x , since it consists of d line segments each with slope $(\tilde{\mathbf{w}}^{\downarrow})_i$ for $i = 1, \dots, d$ which

¹ It satisfies the condition $\text{mana}(\rho_1 \otimes \rho_2) = \text{mana}(\rho_1) + \text{mana}(\rho_2)$ which is practical in distillation settings.

is non-increasing by definition. Finally, the points on the interior of the Lorenz curve connecting line segments of different slopes are called *elbows*.

A vector \mathbf{y} is said to \mathbf{d} -majorise another vector \mathbf{x} if and only if the Lorenz curve $L_{\mathbf{y}|\mathbf{d}}$ lies above Lorenz curve $L_{\mathbf{x}|\mathbf{d}}$, thus reducing \mathbf{d} -majorisation into a finite set of inequalities.

Theorem 4. *Let $\mathbf{x}, \mathbf{y}, \mathbf{d} \in \mathbb{R}^d$, such that the components of \mathbf{d} are positive. Then, $\mathbf{x} \prec_{\mathbf{d}} \mathbf{y}$ if and only if $L_{\mathbf{x}|\mathbf{d}}(x) \leq L_{\mathbf{y}|\mathbf{d}}(x)$ for all $x \in [0, 1]$ with strict equality at $x = 1$.*

A restatement of the theorem including more equivalent conditions and a proof are provided in Appendix B.

III. STOCHASTIC STRUCTURE OF MAGIC THEORIES

A. Magic fragments

Equipped with the definitions of the Wigner distribution in odd prime dimensions, we can formally recast the maximal magic theory \mathcal{R}_{\max} into a stochasticity setting. The free states correspond to proper probability distributions

$$\mathcal{F}_{\max} := \{\rho : W_{\rho}(\mathbf{x}) \geq 0 \text{ for all } \mathbf{x} \in \mathcal{P}_d\} \quad (18)$$

The free operations should send the set of free states \mathcal{F}_{\max} into itself and completely preserve the non-negativity of the states, in the sense that $\mathcal{E} \in \mathcal{O}_{\max}$ iff $(\mathbb{1}_d \otimes \mathcal{E})\sigma \in \text{STAB}$ for all odd prime dimensions d whenever $\sigma \in \mathcal{F}_{\max}$. It is shown by Wang *et al.* [12] that \mathcal{O}_{\max} coincides with the set of operations \mathcal{E} with stochastic Wigner distributions,

$$\mathcal{O}_{\max} = \{\mathcal{E} : W_{\mathcal{E}}(\mathbf{y}|\mathbf{x}) \geq 0 \text{ for all } \mathbf{x}, \mathbf{y} \in \mathcal{P}_d\}. \quad (19)$$

Every established magic theory $\mathcal{R} = (\mathcal{F}, \mathcal{O})$ is a sub-theory of \mathcal{R}_{\max} in the sense that $\mathcal{F} \subseteq \mathcal{F}_{\max}$ and $\mathcal{O} \subseteq \mathcal{O}_{\max}$, and as such it falls under this new stochasticity setting. For technical simplicity in what follows we assume that \mathcal{F} is a closed set, and note that \mathcal{F}_{\max} is itself a closed set, since it is specified by a finite set of linear constraints of the form $\text{tr}[L\rho] \geq 0$ with operators $L \in \mathcal{B}(\mathcal{H})$ ensuring that the state is positive and Wigner-positive.

Given this context we now define the following key notion, that is central to our analysis.

Definition 5 (σ -fragment). *Given a resource theory of magic $\mathcal{R} = (\mathcal{F}, \mathcal{O})$, the σ -fragment of \mathcal{R} is the resource theory $\mathcal{R}_{\sigma} = (\mathcal{F}_{\sigma}, \mathcal{O}_{\sigma})$, where $\mathcal{F}_{\sigma} = \{\sigma\}$ and the free operations are restricted to the ones that leave σ invariant,*

$$\mathcal{O}_{\sigma} := \{\mathcal{E} \in \mathcal{O} : \mathcal{E}(\sigma) = \sigma\}. \quad (20)$$

Note that $\mathcal{F} = \bigcup_{\sigma \in \mathcal{F}} \mathcal{F}_{\sigma}$ trivially. With this basic notion defined, we now show that this union holds for free operations as well and therefore that any resource theory

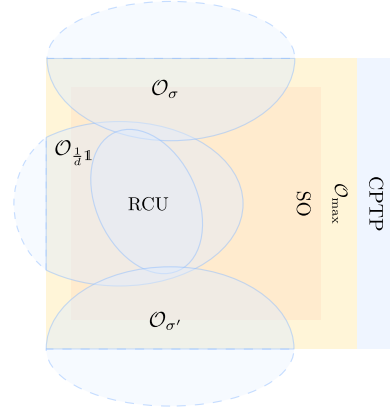


FIG. 2. Decomposition of a magic theory \mathcal{R} into σ -fragments. Examples of magic theories (SO: Stabilizer operations, \mathcal{O}_{\max} : Completely positive-Wigner-preserving operations, RCU: Random Clifford Unitaries – subclass of SO) involve operations denoted by the two yellow regions, with every other established magic theory between them. We introduce σ -fragments \mathcal{O}_{σ} defined for all free states σ that cover \mathcal{O}_{\max} . Each \mathcal{O}_{σ} is extendible to a set of stochastic maps outside the CPTP operations. Within each σ -fragment \mathbf{d} -majorisation can be used allowing for an intricate, but tractable approach towards the study of magic state interconversion.

of magic can be faithfully subdivided into σ -fragments, in such a way that any problem of interconversion in the parent magic theory \mathcal{R} can be analysed across the different fragments.

Theorem 6. *Let $\mathcal{R} = (\mathcal{F}, \mathcal{O})$ be a theory of magic. Every operation in \mathcal{O} leaves at least one free state invariant,*

$$\mathcal{O} = \bigcup_{\sigma \in \mathcal{F}} \mathcal{O}_{\sigma}. \quad (21)$$

Therefore, $\rho \rightarrow \tau$ in \mathcal{R} if and only if $\rho \rightarrow \tau$ in a σ -fragment of \mathcal{R} .

Proof. Suppose \mathcal{E} is in a σ -fragment \mathcal{O}_{σ} . Then it is also in \mathcal{O} , hence $\bigcup_{\sigma \in \mathcal{F}} \mathcal{O}_{\sigma} \subseteq \mathcal{O}$.

Conversely, suppose \mathcal{E} is in \mathcal{O} . The free states are a closed set that is mapped one-to-one to a closed subset \mathcal{S} of the $(d^2 - 1)$ -dimensional probability simplex. \mathcal{S} is convex, since any combination of free states is also free and the Wigner distribution is linear. Therefore, \mathcal{S} is convex and compact as a closed convex subset of the bounded compact probability simplex.

We can now view $W_{\mathcal{E}}$ as a stochastic, continuous mapping from \mathcal{S} to itself, thus Brouwer's fixed point theorem [CITE] implies that there exists a probability distribution $d_{\mathbf{z}}$ for some $\mathbf{z} \in \mathcal{P}_d$ that is a fixed point of $W_{\mathcal{E}}$. This corresponds to a free state $\sigma := \sum_{\mathbf{z} \in \mathcal{P}_d} d_{\mathbf{z}} A_{\mathbf{z}} \in \mathcal{F}$. Therefore $\mathcal{E} \in \mathcal{O}_{\sigma}$, and so $\mathcal{O} = \bigcup_{\sigma \in \mathcal{F}} \mathcal{O}_{\sigma}$.

The state interconversion result follows immediately. \square

The zoo of all magic operation classes is summarised in Fig. (2). Completely positive-Wigner-preserving operations [12] form the maximal operation class \mathcal{O}_{\max} . This theory of magic is thus covered by σ -fragments exactly and every other magic theory is contained within this cover. In particular, the stabiliser operations SO are contained within \mathcal{O}_{\max} .

Every quantum circuit aiming at a given magic state conversion $\rho \rightarrow \tau$ possesses certain symmetries according to Theorem 6 that allow us to study the conversion within only certain σ -fragments that reflect these symmetries. As a simple example, the dephasing channel

$$\Delta(\rho) = \sum_{k \in \mathbb{Z}_d} |k\rangle\langle k| \rho |k\rangle\langle k| \quad (22)$$

removes coherent phases in the computational basis and therefore leaves exactly all mixtures of computational basis states invariant. A noise channel consisting only of such dephasing operations can be fully analysed in the σ -fragments for the pure computational basis states σ . [Have proven in appendix that we need to consider *only* the pure stabiliser fragments - not mixed ones]

The subdivision of magic theories into σ -fragments is also powerful because the pre-order $\prec_{\mathcal{R}_\sigma}$ of every σ -fragment is described by well-behaved majorisation tools, as we establish in the rest of this section.

B. Majorisation of quasi-probabilities in σ -fragments

We can approach any magic theory through a thermodynamic lens, and in doing so we are provided with valuable insights on the structure of the theory. Firstly, any free state, for example a stabiliser state, can be viewed as a thermal state γ_β . Without loss of generality we can always write a quantum state σ as a thermal state, $\sigma = \gamma_\beta := \frac{1}{\mathcal{Z}_\beta} e^{-\beta H}$ for some $\beta \geq 0$ and Hamiltonian H (either effective or actual)². We can also view the set of free operations \mathcal{O}_σ as a counterpart of thermal operations, in the sense that any operation \mathcal{E} in \mathcal{O}_σ preserves state σ .

It is then apparent that the pre-order $\prec_{\mathcal{R}_\sigma}$ between the operations in the σ -fragment follows the rules of d -majorisation as outlined in Section IID. In particular, the pre-order $\prec_{\mathcal{R}_\sigma}$ of the σ -fragment $\mathcal{R}_\sigma = (\mathcal{F}, \mathcal{O}_\sigma)$ between d -dimensional states corresponds to the majorisation pre-order \prec_{W_σ} between their d^2 -dimensional Wigner distributions. For simplicity we shall merge the notation into \prec_σ , as there is little risk of confusion.

Note that the Wigner components of an n -copy state $\rho^{\otimes n}$ can be calculated directly from W_ρ by convolution

of the distribution with itself,

$$W_{\rho^{\otimes n}} = W_\rho^{\otimes n}, \quad (23)$$

where \otimes can be interpreted as the usual Kronecker product in the last expression. We may use the vector notation $\mathbf{w}(\rho) = (w(\rho)_i)_{i=1, \dots, d^2}$ for the Wigner distribution, in which case we can express the Kronecker product as

$$W_\rho \otimes W_\rho = (w(\rho)_i w(\rho)_j)_{i,j=1, \dots, d^2}. \quad (24)$$

The correspondence between the phase space and pure vector representations of the Wigner distribution is discussed more in Appendix F 1.

Furthermore, we restrict our analysis to σ -fragments, where σ is full-rank. This is justified as majorisation is continuous between fragments, in the sense that the pre-orders \prec_σ and $\prec_{\sigma'}$ are equivalent for states σ and σ' which are ϵ -close. [CHECK] Therefore, majorisation analysis is robust under imperfections in the experimental implementation of quantum operations.

In particular, we can always use common noise effects to approximate any σ -fragment where σ is not full-rank by the σ' -fragment where σ' is a noisy approximation of σ . For example, inducing depolarising noise, we can write $\sigma' = (1 - \epsilon)\sigma + \epsilon \frac{1}{d} \mathbb{1}$, for some infinitesimal $\epsilon > 0$, so that σ' is full-rank and arbitrarily close to σ . Important examples of such σ -fragments include pure stabiliser states which are rank-1, e.g. the zero state depicted in Fig. (1(b)). Operations in such fragments include important stabiliser operations like the replacement channel, $\mathcal{E}(\rho) = \sigma$ for all states ρ .

Theorem 7. *Let $\mathcal{R} = (\mathcal{F}, \mathcal{O})$ be a theory of magic. If $\rho \rightarrow \tau$ in \mathcal{R} then $W_\tau \prec_\sigma W_\rho$ within at least one σ -fragment.*

Proof. Suppose we can convert ρ into τ in the magic theory. Thus there is some \mathcal{O}_σ and some $\mathcal{E} \in \mathcal{O}_\sigma$ such that $\mathcal{E}(\rho) = \tau$, and $\mathcal{E}(\sigma) = \sigma$. Therefore, the Wigner distribution of this free operation satisfies $W_\mathcal{E} \in \text{Sd}^2(W_\sigma)$ and $W_\mathcal{E} W_\rho = W_\tau$. Since σ is full-rank and free, its Wigner distribution is strictly positive in all components, so it directly follows from Definition 2 that $W_\tau \prec_\sigma W_\rho$. \square

The result of Theorem 7 can be understood as an extension of the idea of a magic monotone, where we replace $\mathcal{M}(\tau) \leq \mathcal{M}(\rho)$ with $W_\tau \prec_\sigma W_\rho$. The physical difference between the two expressions is that the majorisation ordering occurs in a specific σ -fragment. Therefore, majorisation constraints can be used to place upper bounds on magic state distillation in a way that allows one to incorporate the physics of the allowed operations – specifically, it enables us to bound how much magic can be distilled via quantum operations that, for example, preserve the equilibrium state of the system, or via operations that are symmetric about the Z -axis of the Bloch sphere. We discuss distillation upper bounds in detail in Sections IV and V

This approach can also provide *lower bounds* on distillation, however now more structure about the specific

² Technicalities arise for the case where σ is not full rank, but this can be still described via $\beta \rightarrow 0$ limiting process.

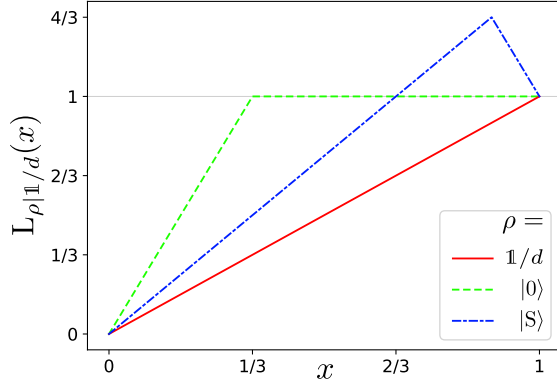


FIG. 3. **Quasi-probability Lorenz curve comparison.** The Lorenz curves are constructed by the Wigner distributions illustrated in Fig. (1) in the unital fragment $\mathcal{O}_{\mathbb{1}/d}$ for $d = 3$. The maximally mixed state curve is simply the line connecting $(0,0)$ and $(1,1)$. There is no operation in $\mathcal{O}_{\mathbb{1}/d}$ that can convert $|S\rangle$ to $|0\rangle$, as their Lorenz curves intersect.

free operations must be included. We briefly discuss distillation lower bounds in [REF].

[Mention relative majorisation as a possibility]

C. Lorenz curves of quasi-probabilities in σ -fragments

It is straightforward to construct Lorenz curves for Wigner distributions in any σ -fragment. As we have seen in Theorem 7, for any full-rank free state σ we have that W_σ is a strictly positive full-rank probability distribution, and so one can define a corresponding notion of d -majorisation on *quasi*-distributions. We write $L_{\rho|\sigma}(x)$ for the Lorenz curve of W_ρ with respect to W_σ . The vector of ratios $\tilde{\mathbf{w}}(\rho|\sigma)$ used to construct the curve is

$$\tilde{w}(\rho|\sigma)_i := \frac{w(\rho)_i}{w(\sigma)_i}, \quad (25)$$

and is called the *Gibbs-rescaled* distribution of ρ with respect to σ .

An example of comparison between different Lorenz curves is provided in Fig. (3). The curves in the figure are constructed in the *unital fragment*, i.e. the σ -fragment defined by the maximally mixed state $\sigma = \frac{1}{d}\mathbb{1}$ whose Wigner distribution is the uniform probability distribution.

Normalisation of the Wigner distribution [REF] ensures that for all quantum states ρ , $L_{\rho|\sigma}(x) \geq 0$ and $L_{\rho|\sigma}(1) = 1$. We stress that $0 \leq L_{\rho|\sigma}(x) \leq 1$ for all $x \in [0, 1]$ if and only if ρ is a positive Wigner state. As a consequence, checking whether a magic state conversion of the form

$$\rho \xrightarrow{\mathcal{E} \in \mathcal{O}_\sigma} \tau \quad (26)$$

is not possible, reduces to the set of constraints

$$L_{\rho|\sigma}(x) \geq L_{\tau|\sigma}(x), \quad x \in [0, 1], \quad (27)$$

due to Theorems 4 and 7.

We can refine the number of independent constraints stemming from this inequality. In fact, there are only as many independent constraints as there are elbows in the Lorenz curve of the target state as shown in Appendix D 4. However, in principle any one inequality defined at some location x provides a valid constraint leading to some upper distillation bound, and the more locations one considers, the stricter the bound is.

The area $\mathcal{A}_\sigma(\rho)$ between the curve $L_{\rho|\sigma}$ and the line $y = 1$ is a resource monotone in the σ -fragment. This is clear because for any state conversion like Eq. (26), the Lorenz curve $L_{\tau|\sigma}$ is lower than $L_{\rho|\sigma}$, hence $\mathcal{A}_\sigma(\mathcal{E}(\rho)) \leq \mathcal{A}_\sigma(\rho)$. The exact form of the monotone is given in Appendix D 3. [should we include the area? Have a look at Appendix D 3]

We can now show that the magic monotone of mana, and equivalently the sum-negativity of a quantum state, can be viewed as a feature of majorisation, and independent of the particular σ -fragment one works in. In fact, our current setting makes it apparent that mana provides a weaker condition than majorisation for all magic state interconversions.

We first show that the Lorenz curve maximum of state ρ is independent of the σ -fragment and directly related to its sum-negativity.

Lemma 8. *Given a quantum state ρ , the maximum of its Lorenz curve $L_{\rho|\sigma}$ is independent of the σ -fragment and is equal to $1 + \text{sn}(\rho)$.*

Proof. We may use the vector notation of the Wigner distributions $\mathbf{w}(\rho)$ and $\mathbf{w}(\sigma)$. We choose the component indexing so that the Gibbs-rescaled distribution

$$\tilde{\mathbf{w}}(\rho|\sigma) := \left(\frac{w(\rho)_1}{w(\sigma)_1}, \dots, \frac{w(\rho)_{d^2}}{w(\sigma)_{d^2}} \right)^T, \quad (28)$$

is sorted, $\tilde{\mathbf{w}} = \tilde{\mathbf{w}}^\downarrow$. Note that all components of $\mathbf{w}(\sigma)$ are positive, so $\tilde{w}_i \geq 0$ if and only if $w(\rho)_i \geq 0$ for any $i = 1, \dots, d^2$.

Let i_\star be the index of the smallest non-negative component of $\tilde{\mathbf{w}}^\downarrow$. Then, $w(\rho)_i < 0$ if and only if $i > i_\star$, so the maximum of Lorenz curve $L_{\rho|\sigma}(x)$ takes the value

$$L_{\rho|\sigma}(x_{i_\star}) = \sum_{i=1}^{i_\star} w(\rho)_i, \quad (29)$$

and is achieved at

$$x_{i_\star} := \sum_{i=1}^{i_\star} w(\sigma)_i. \quad (30)$$

The location of the maximum ($x = x_{i_\star}$) varies from fragment to fragment, but its value is independent of σ ,

$$L_{\rho|\sigma}(x_{i_\star}) = \sum_{\mathbf{x}: W_\rho(\mathbf{x}) \geq 0} W_\rho(\mathbf{x}) = 1 + \text{sn}(\rho). \quad (31)$$

□

We can therefore view mana as just one feature of the Lorenz curve, namely its maximum value. Conversely, it is now clear that the maximum of the Lorenz curve acts as a valid magic monotone. This implies the following result.

Theorem 9. *Given a magic state conversion $\rho \rightarrow \tau$, the majorisation condition is stronger than the mana condition in every σ -fragment.*

Proof. The maximum of the Lorenz curve of a state ρ is independent of the σ -fragment due to Lemma 8, and can be expressed as an increasing function of mana,

$$\max_{x \in [0,1]} L_{\rho|\sigma}(x) = 1 + \text{sn}(\rho) = \frac{1}{2} \left(1 + e^{\text{mana}(\rho)} \right). \quad (32)$$

Therefore, the majorisation bound

$$L_{\rho|\sigma}(x) \geq L_{\tau|\sigma}(x), \quad x \in [0,1] \quad (33)$$

implies the order $\max_{x \in [0,1]} L_{\rho|\sigma}(x) \geq \max_{x \in [0,1]} L_{\tau|\sigma}(x)$, hence the mana condition $\text{mana}(\rho) \geq \text{mana}(\tau)$. \square

In sections Sections IV and V, we study majorisation constraints on magic distillation arising in different fragments. For this reason, we define the n -copy, ϵ -noisy Strange state,

$$\rho_S(\epsilon)^{\otimes n} := \left[(1 - \epsilon) |S\rangle \langle S| + \epsilon \frac{1}{3} \mathbb{1} \right]^{\otimes n}, \quad (34)$$

in \mathcal{O}_σ , where the pure magic state $|S\rangle$ is induced with depolarising noise. Its Wigner distribution is visualised in Fig. (1(d)).

At $\epsilon = 0$, the Strange state is a qutrit magic state of maximal sum-negativity / mana [11] and therefore acts as an ideal distillation target, analogous to a Bell state in bipartite entanglement theory. The Strange state is exceptionally symmetric under Clifford transformations and as a result there exists a *twirling* protocol, discussed in detail in [31, 32], allowing for the conversion of any noisy magic state to the form of Eq. (34) via Cliffords.

In Fig. (4), we illustrate the Lorenz curve $L_{\rho_S|\sigma}^3$ of pure and noisy n -copy Strange states in some thermal σ -fragment. Due to Lemma 8, it is clear that the curves peak at $1 + \text{sn}(\rho_S(\epsilon)^{\otimes n})$.

IV. MAGIC IN THE UNITAL FRAGMENT

The unital fragment encompasses the circuits which preserve the maximally mixed state ($1/d$) and so it includes many important families of circuits.

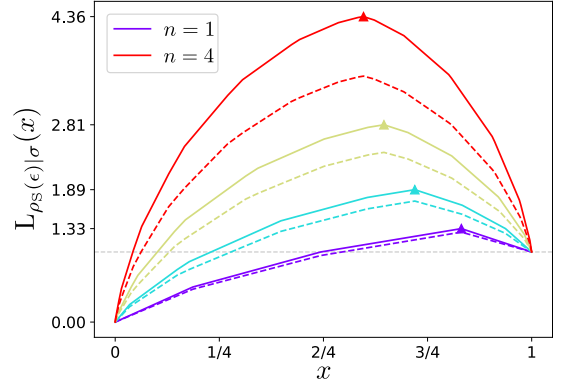


FIG. 4. **Strange state Lorenz curves.** Lorenz curves of $\rho_S(\epsilon)^{\otimes n}$ for $n = 1, 2, 3, 4$ in the σ -fragment where $\sigma = e^{-\beta H} / \mathcal{Z}_\beta$, with Hamiltonian spectrum $(0, 1, 2)$ over the computational basis and $\beta = 0.5$. Solid lines represent pure Strange states ($\epsilon = 0$); dashed lines represent noisy Strange states ($\epsilon = 0.1$).

MSD circuits in principle consist of bulk sequences of random Clifford unitaries (RCU)[CITE], depicted in Fig. (2). Operations in RCU can be expressed as

$$\mathcal{E}(\rho) = \sum_i p_i U_i \rho U_i^\dagger, \quad U_i \in \mathcal{C}_d. \quad (35)$$

Depending on the symmetries of such operations, a Clifford sequence may belong in other σ -fragments as well. In such a case, the majorisation condition (27) needs to be checked in the σ -fragments that reflect all symmetries of the operation sequence.

In general, noisy circuits are well-described by the unital fragment. To see this, consider incorporating noisy channels in the circuit, for example dephasing channels as in Eq. (22) defined in different bases. This process destroys the circuit symmetries, except for the invariance of the maximally mixed state. Dephasing and bit-flip error channels are examples of the many error-inducing channels that respect the unital symmetry. [\[Expand on significance of unital fragment\]](#)

We now consider the task of purifying n copies of a noisy Strange state $\rho_S(\epsilon)^{\otimes n}$ as given in Eq. (34) into a smaller number of copies n' of a less noisy strange state $\rho_S(\epsilon')^{\otimes n'}$, with $\epsilon' < \epsilon$ and $n' \leq n$,

$$\rho_S(\epsilon)^{\otimes n} \rightarrow \rho_S(\epsilon')^{\otimes n'} \otimes \left(\frac{1}{3} \mathbb{1} \right)^{\otimes (n-n')}, \quad (36)$$

where all copies $n, n', n - n'$ are even. Since the state $1/3$ is free, tensoring in copies of it does not affect the distillation process.

The Lorenz curve $L_{\rho_S|1/3}$, for some general noise parameter ϵ and number of copies n , is defined at 9^n points between 0 and 1. The exact expressions for the coordinates of these points can take 8 different forms, depending on whether the noise level [\[noise level / error rate\]](#) ϵ is greater or less than $\frac{3}{7}$, the parity of the number of

³ Hereinafter, we may drop parameters ϵ and n for clarity

copies n is even or odd and the location relative to the curve peak is on the left hand side (LHS – including the curve peak) or right hand side (RHS) of the curve peak. The full details for the construction of all Lorenz curve forms are provided in Appendix E.

[In the following, should we talk about the coordinates of all elbows and then say that we are gonna use only the first elbow, or should we leave full-elbow analysis for the appendix and only talk about first elbow here?]

Here we focus on the case of Lorenz curves with even copies n, n' and low noise levels ($\epsilon' < \epsilon \leq 3/7$). The part of the curve on the LHS of its peak is constructed by the positive components of the state Wigner distribution that take values w_i , with associated multiplicities m_i ,

$$w_i = \left(\frac{1}{6} - \frac{1}{18}\epsilon\right)^{2i} \left(-\frac{1}{3} + \frac{4}{9}\epsilon\right)^{n-2i}, \quad (37)$$

$$m_i = 8^{2i} \binom{n}{2i}, \quad (38)$$

where $i = 0, \dots, \frac{n}{2}$. It is readily seen that $L_{\rho_S|1/3}(x)$ has elbows on the LHS of its peak located at points $x = x_i$, where

$$x_i = \Phi_+\left(2i; n, \frac{8}{9}\right), \quad (39)$$

with $i = 0, 1, \dots, n$ and where $\Phi_+(i, n, p)$ is an even-power cumulant distribution function given by

$$\Phi_+(2i; n, p) := \sum_{\ell=0}^i \binom{n}{2\ell} p^{2\ell} (1-p)^{n-2\ell}. \quad (40)$$

We provide technical details on the function Φ_+ in the appendices.

The value of the Lorenz curve at an elbow is given by

$$L_{\rho|1/3}(x_i) \equiv L_i = \left(\frac{5}{3} - \frac{8}{9}\epsilon\right)^n \Phi_+\left(2i; n, \frac{12-4\epsilon}{15-8\epsilon}\right), \quad (41)$$

and we define $(x_{-1}, L_{-1}) = (0, 0)$.

Lorenz curve comparison provides 9^n necessary constraints for the process Eq. (36). Theorem 18 shows that they can be reduced to n' independent constraints. However, in principle we can pick any location $x_\star \in [0, 1]$ at which the Lorenz curves are defined and we then get a necessary distillation constraint,

$$L_{\rho_S(n, \epsilon, 0)|1/3}(x_\star) \geq L_{\rho_S(n', \epsilon', n-n')|1/3}(x_\star). \quad (42)$$

We choose $x_\star = x_0 = \Phi_+(0, n, \frac{8}{9}) = \frac{1}{9^n}$, the location of the first elbow of the initial state $\rho_S(n, \epsilon, 0)$. The initial state Lorenz curve coordinate at this location is

$$L_\star = L_0 = \left(\frac{5}{3} - \frac{8}{9}\epsilon\right)^n \Phi_+\left(0; n, \frac{12-4\epsilon}{15-8\epsilon}\right) = \left(\frac{3-4\epsilon}{9}\right)^n \quad (43)$$

The first elbow of the target state is located at $x'_0 = \frac{1}{9^{n'}} > x_\star$, so in order to find the target state Lorenz curve

coordinate L'_\star at location x_\star , we can interpolate between the origin and the first target state elbow,

$$\begin{aligned} L'_\star &= \frac{x_\star}{x'_0} L'_0 = \frac{1}{9^{n-n'}} \left(\frac{5}{3} - \frac{8}{9}\epsilon'\right)^{n'} \Phi_+\left(0; n', \frac{12-4\epsilon'}{15-8\epsilon'}\right) \\ &= \frac{(3-4\epsilon')^{n'}}{9^n}. \end{aligned} \quad (44)$$

The process in Eq. (36) is therefore possible only if $L_\star \geq L'_\star$ or equivalently,

$$\frac{n}{n'} \geq \frac{\ln(3-4\epsilon')}{\ln(3-4\epsilon)}, \text{ or} \quad (45)$$

$$\epsilon \leq \frac{3}{4} - \frac{1}{4}(3-4\epsilon')^{\frac{n'}{n}}. \quad (46)$$

As an example, the process with $(n, n') = (12, 4)$ and $\epsilon' = 0.05$ yields $\epsilon \leq 0.398$. Numerical comparison of the curves with varying ϵ suggests that we require $\epsilon \leq 0.373$ for the process to be possible.

Moreover, for the problem of distilling pure magic states where $\epsilon' = 0$ we obtain an upper bound on the unital fragment distillation rate $R = n'/n$ given by

$$R \leq 1 + \frac{\ln(1 - \frac{4}{3}\epsilon)}{\ln 3}, \quad (47)$$

which holds for all finite n . In comparison, the bound from mana gives

$$R_{\text{mana}} \leq 1 + \frac{\ln(1 - \frac{8}{15}\epsilon)}{\ln \frac{5}{3}}, \quad (48)$$

a threshold which is looser than the one provided by majorisation for all $\epsilon > 0$ as illustrated in Fig. (5). For example, for $\epsilon = 0.4$ we obtain $R \leq 0.31$ while the mana bound gives $R_{\text{mana}} \leq 0.53$.

V. MAGIC BOUNDS IN ARBITRARY STABILISER FRAGMENTS

We now generalise the approach taken in the previous section and consider bounds on magic distillation for an arbitrary stabiliser fragment, \mathcal{R}_σ where σ is any quantum state $\sigma \in \text{STAB}$. In other words, we consider those bounds on distillation that apply when the free operations have σ as a fixed point.

These bounds can be interpreted in two different ways: on one hand they can be viewed as a family of upper bounds parameterized by a stabiliser state σ , on the other we can associate σ to actual hardware limitations or to biased noise models in which it is an equilibrium state of some kind. Without loss of generality, we can always write a stabiliser state σ as a Gibbs state $\sigma = \gamma_\beta := \frac{1}{\mathcal{Z}_\beta} e^{-\beta H}$ for some $\beta \geq 0$ and Hamiltonian H as discussed in Section IIIB.

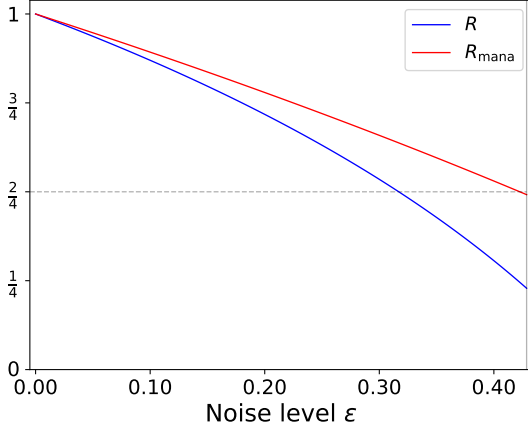


FIG. 5. **Distillation rate in unital fragment.** Distillation rates obtained by mana and majorisation are plotted for $\epsilon' = 0$, up to noise levels $\epsilon_{\max} = \frac{3}{7}$. Majorisation consistently provides stricter rates as proven in Theorem 9. [add thauma?]

We focus on the distillation process

$$\rho_S(\epsilon)^{\otimes n} \longrightarrow \rho_S(\epsilon')^{\otimes n'} \otimes \sigma^{\otimes (n-n')}, \quad (49)$$

where the noisy Strange state is given in Eq. (34) and we highlight again that any magic state can be transformed to this form via Clifford operations. Notice that tensoring in copies of γ_β does not affect the process, since the state is free. In the following result, all copies $n, n', n - n'$ are even with $n > n'$, and $\epsilon' = 0$, but the bounds are easily generalised to odd numbers of copies and ϵ' such that $0 < \epsilon' < \epsilon$. Finally, we assume a range of initial noise levels, $\epsilon \leq 3/7$ [higher than all threshold values in literature CITE], so that the largest Wigner component of the 1-copy noisy magic state is negative.

Given this context, we now provide the following result, which depends on the free energy F_β of state σ ,

$$F_\beta := \text{tr}[H\sigma] - \beta^{-1}S(\sigma) = -\beta^{-1} \log \mathcal{Z}_\beta, \quad (50)$$

where the von Neumann entropy is $S(\sigma) := -\text{tr}[\sigma \log \sigma]$.

Theorem 10. *Let σ be a qutrit stabiliser state, given by $\sigma = \frac{1}{\mathcal{Z}} e^{-\beta H}$ with H having eigenvalues $E_0 \leq E_1 \leq E_2$, and where β is an (effective) inverse temperature for the state. We define*

$$\beta_\star := \frac{1}{E_2 - E_0} \ln 2 \quad (51)$$

and for $\beta \leq \beta_\star$, we define a threshold error rate

$$\epsilon_\star(\beta) := 3 - \frac{18}{8 - e^{(E_2 - E_0)\beta}}. \quad (52)$$

Then any distillation rate $R(\epsilon, \beta)$ in the σ -fragment of a magic theory is bounded as follows:

If $\beta \leq \beta_\star$ and $\epsilon \leq \epsilon_\star$,

$$R(\epsilon, \beta) \leq 1 + \frac{\ln(1 - \frac{4}{3}\epsilon)}{\beta(E_0 - F_\beta)}. \quad (53)$$

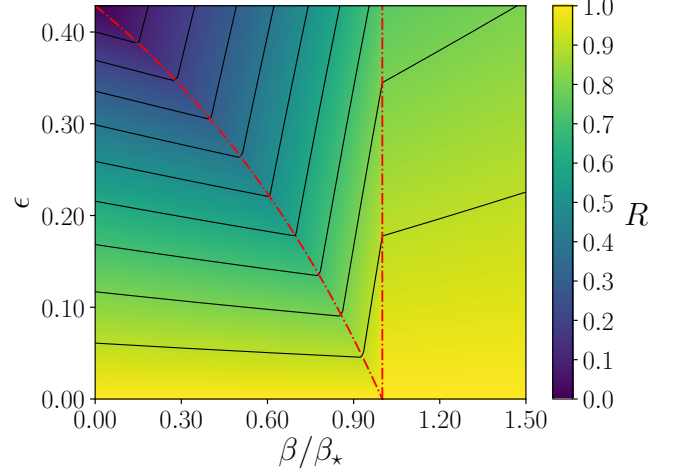


FIG. 6. **Bounds on magic distillation rates $R(\epsilon, \beta)$ within stabiliser fragments.** The vertical dashed line is the ‘Landauer-like’ temperature threshold β_\star and the diagonal dashed curve corresponds to the noise threshold $\epsilon_\star(\beta)$ at every $\beta \leq \beta_\star$. The unital fragment corresponds to the $\beta = 0$ line.

If $\beta \leq \beta_\star$ and $\epsilon_\star < \epsilon$, then

$$R(\epsilon, \beta) \leq 1 + \frac{\ln(1 - \frac{1}{3}\epsilon) - (E_2 - E_0)(\beta_\star - \beta)}{\beta(E_0 - F_\beta)}. \quad (54)$$

Otherwise if $\beta > \beta_\star$,

$$R(\epsilon, \beta) \leq 1 + \frac{\ln(1 - \frac{1}{3}\epsilon)}{-\ln 2 + \beta(E_2 - F_\beta)}. \quad (55)$$

A few comments can be given on this result. Firstly, the specific numerical factors in ϵ_\star are a result of our choice of magic state. Secondly, these bounds are derived based on analysis of only a part of the Lorenz curves and can be improved via a finer analysis. This is apparent by simple numerical calculations on the entirety of the curves. The existing bounds follow by considering the dominant terms in the rescaled Wigner distribution and do not, for example, take into account the Lorenz curve’s peak structure. Finally, it is striking that we obtain a Landauer-like condition with a characteristic temperature $kT_\star \ln 2 = E_2 - E_0$, where $kT_\star = \beta_\star^{-1}$. It is unclear whether this points to a generic feature that can be directly related to fundamental thermodynamic relations, such as the erasure cost of a single bit being $kT \ln 2$. Given that we work with qutrits this seems surprising, but does deserve further study.

Proof. We provide an outline of the proof here, with all technical details included in [the appendix].

The Wigner distribution of ϵ -noisy Strange state can be written as

$$W_{\rho_S(\epsilon)}(\mathbf{x}) = (1 - \epsilon)W_{|S\rangle\langle S|}(\mathbf{x}) + \epsilon W_{\frac{1}{3}\mathbf{1}}(\mathbf{x}), \quad (56)$$

so we get positive components

$$u(\epsilon) = \frac{1}{6} - \frac{1}{18}\epsilon \quad (57)$$

at the 8 phase space points $\mathbf{x} \in \mathcal{P}_d[3] \setminus \{\mathbf{0}\}$ and a negative component

$$-v(\epsilon) = -\left(\frac{1}{3} - \frac{4}{9}\epsilon\right) \quad (58)$$

at the origin $\mathbf{x} = \mathbf{0}$.

We are interested in studying Strange state distillation in general thermal fragments $\mathcal{O}_{\gamma_\beta}$ where the thermal state is

$$\gamma_\beta := \frac{e^{-\beta H}}{\mathcal{Z}_\beta}, \quad (59)$$

with partition function $\mathcal{Z}_\beta = \text{tr}[e^{-\beta H}]$, defined in terms of some Hamiltonian H and temperature β^{-1} .

To this end, consider the general qutrit Hamiltonian

$$H = E_0 |0\rangle\langle 0| + E_1 |1\rangle\langle 1| + E_2 |2\rangle\langle 2|, \quad (60)$$

where the energy eigenbasis coincides with the computational basis. We have that energy eigenvalues $E_i \geq 0$, but impose no order between them. We denote the highest energy eigenvalue as

$$E_{\max} := \max\{E_0, E_1, E_2\}. \quad (61)$$

The thermal state associated with this Hamiltonian is

$$\gamma_\beta = \frac{e^{-\beta H}}{\mathcal{Z}_\beta} = \frac{e^{-\beta \sum_{k=0}^2 E_k |k\rangle\langle k|}}{\mathcal{Z}_\beta} \quad (62)$$

$$= \frac{e^{-\beta E_0}}{\mathcal{Z}_\beta} |0\rangle\langle 0| + \frac{e^{-\beta E_1}}{\mathcal{Z}_\beta} |1\rangle\langle 1| + \frac{e^{-\beta E_2}}{\mathcal{Z}_\beta} |2\rangle\langle 2|, \quad (63)$$

where $\mathcal{Z}_\beta = e^{-\beta E_0} + e^{-\beta E_1} + e^{-\beta E_2}$.

The Wigner distribution of state γ_β can be seen as the ensemble average of the distributions of the computational basis states,

$$W_{\gamma_\beta}(\mathbf{x}) = \sum_{k=0}^2 \frac{e^{-\beta E_k}}{\mathcal{Z}_\beta} W_{|k\rangle\langle k|}(\mathbf{x}) \quad (64)$$

$$= \sum_{k=0}^2 \frac{e^{-\beta E_k}}{\mathcal{Z}_\beta} \delta_{x_0, k} = \frac{e^{-\beta E_{x_0}}}{3\mathcal{Z}_\beta}, \quad (65)$$

for all $\mathbf{x} \in \mathcal{P}_d[3]$. All Wigner components are strictly positive, therefore the pre-order \prec_{γ_β} is well defined.

Let $\sigma = \frac{1}{\mathcal{Z}} e^{-\beta H}$ be a stabiliser state, which is given in its eigen-decomposition as

$$\sigma = \frac{1}{\mathcal{Z}} (e^{-\beta E_0} |\varphi_0\rangle\langle \varphi_0| + e^{-\beta E_1} |\varphi_1\rangle\langle \varphi_1| + e^{-\beta E_2} |\varphi_2\rangle\langle \varphi_2|), \quad (66)$$

where $\{|\varphi_k\rangle\}$ are pure, orthonormal stabiliser states. Therefore there exist a Clifford operation that takes $|\varphi_k\rangle \rightarrow |k\rangle$, with the effect of permuting the negative component of $\rho_S(\epsilon)$ in the discrete phase space.

We are now in a position to consider the Strange state distillation process

$$\rho_S(\epsilon)^{\otimes n} \xrightarrow{\mathcal{E} \in \mathcal{O}_{\gamma_\beta}} \rho_S(\epsilon')^{\otimes n'} \otimes \gamma_\beta^{\otimes (n-n')}, \quad (67)$$

where we are allowed to tensor in copies of the thermal state γ_β , since it is free in fragment $\mathcal{O}_{\gamma_\beta}$.

Our aim is to obtain a distillation bound for Eq. (67) which depends on variables $n, n', \epsilon, \epsilon'$ as well as β . In the analysis that follows, we drop obvious variable dependencies for clarity.

Consider a general ϵ -noisy Strange state ρ_S in a thermal fragment \mathcal{O}_γ . We can construct the Gibbs-rescaled distribution

$$\widetilde{W}_{\rho_S|\gamma}(\mathbf{x}) := \frac{W_{\rho_S}(\mathbf{x})}{W_\gamma(\mathbf{x})}, \quad (68)$$

Notice that the Gibbs-rescaled distribution is multiplicative in general,

$$\begin{aligned} \widetilde{W}_{\rho_A \otimes \rho_B | \gamma_A \otimes \gamma_B}(\mathbf{x}_A \oplus \mathbf{x}_B) &= \frac{W_{\rho_A \otimes \rho_B}(\mathbf{x}_A \oplus \mathbf{x}_B)}{W_{\gamma_A \otimes \gamma_B}(\mathbf{x}_A \oplus \mathbf{x}_B)} = \\ &= \frac{W_{\rho_A}(\mathbf{x}_A) W_{\rho_B}(\mathbf{x}_B)}{W_{\gamma_A}(\mathbf{x}_A) W_{\gamma_B}(\mathbf{x}_B)} = \widetilde{W}_{\rho_A | \gamma_A}(\mathbf{x}_A) \widetilde{W}_{\rho_B | \gamma_B}(\mathbf{x}_B). \end{aligned} \quad (69)$$

We construct component and multiplicity vectors based on the distinct components of the Gibbs-rescaled distribution,

$$\mathbf{m} := (1, 2, 3, 3) \quad (70)$$

$$\mathbf{w}(\rho_S) := (-v, u, u, u) \quad (71)$$

$$\mathbf{w}(\gamma) := \frac{1}{3\mathcal{Z}} (e^{-\beta E_0}, e^{-\beta E_0}, e^{-\beta E_1}, e^{-\beta E_2}) \quad (72)$$

$$\mathbf{w}(\rho_S|\gamma) := 3\mathcal{Z} (-ve^{\beta E_0}, ue^{\beta E_0}, ue^{\beta E_1}, ue^{\beta E_2}) \quad (73)$$

There are four distinct values of the Gibbs-rescaled distribution on the phase space, with the negative Strange component $-v$ corresponding to energy E_0 as illustrated in the phase space split in Fig. (7). The Wigner-component vectors $\mathbf{w}(\rho_S), \mathbf{w}(\gamma)$ are then constructed so as to match the four phase space regions.

For notational convenience, we define the index vector

$$\mathbf{q} := (n - \alpha, i, j, k), \quad (74)$$

where i, j, k are non-negative integers and

$$\alpha := i + j + k \leq n.$$

The Wigner components of the n -copy distributions are now readily provided by Lemma 22. Firstly, the n -copy multiplicity is given according to Eq. (F3) by

$$m_{ijk} := M_{\mathbf{q}} = \frac{n!}{i!j!k!(n-\alpha)!} 2^i 3^j 3^k. \quad (75)$$

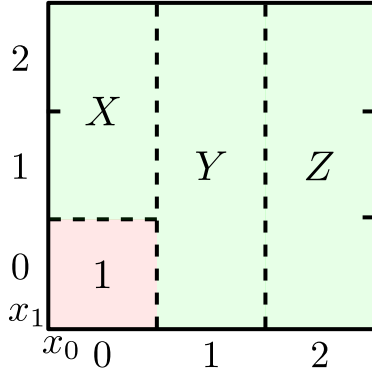


FIG. 7. **Qutrit phase space regions with different Gibbs-rescaled values.** The Gibbs-rescaled distribution attains a unique value in each region 1, X, Y, Z, given by Eq. (73) respectively.

The distribution values that correspond to the same index \mathbf{q} are given according to Eq. (F4) by

$$w(\rho_S)_{ijk} = (-v)^{n-\alpha} u^\alpha \quad (76)$$

$$w(\gamma)_{ijk} = (3\mathcal{Z})^{-n} e^{-\beta(n-\alpha)E_0} e^{-\beta(iE_0+jE_1+kE_2)} \quad (77)$$

$$w(\rho_S|\gamma)_{ijk} = (3\mathcal{Z})^n (-v)^{n-\alpha} u^\alpha e^{\beta(n-\alpha)E_0} e^{\beta(iE_0+jE_1+kE_2)} \quad (78)$$

In order to construct the n -copy Lorenz curve of a Strange state in a thermal fragment we need to sort the components of the n -copy Gibbs-rescaled distribution, $w(\rho_S|\gamma)_{ijk}$ in decreasing order. In particular, to find the coordinates of the first elbow (x_0, L_0) , we need to evaluate the greatest component,

$$w(\rho_S|\gamma)_{\max} := (3\mathcal{Z})^n \max_{i,j,k} \left\{ (-v)^{n-\alpha} u^\alpha e^{\beta(n-\alpha)E_0} e^{\beta(iE_0+jE_1+kE_2)} \right\}, \quad (79)$$

where $0 \leq i, j, k \leq n$ and $\alpha := i + j + k \leq n$. Notice that for $0 \leq \epsilon \leq 3/7$, we have $v \geq u$. We assume that n is even, so we then need the sum $\alpha = i + j + k$ to be even, so that the expression is positive.

Given an even value α for the sum, the term $v^{n-\alpha} u^\alpha e^{-\beta(n-\alpha)E_0}$ is fixed, so the expression is maximised by setting the coefficient of the highest energy E_{\max} equal to α . Hence, we have

$$\begin{aligned} w(\rho_S|\gamma)_{\max} &= (3\mathcal{Z})^n \max_{\substack{\alpha=0,2, \\ \dots, n-2, n}} \left\{ v^{n-\alpha} u^\alpha e^{\beta(n-\alpha)E_0} e^{\beta\alpha E_{\max}} \right\} = \\ &= (3\mathcal{Z})^n v^n e^{n\beta E_0} \max_{\substack{\alpha=0,2, \\ \dots, n-2, n}} \left\{ \left(\frac{u}{v} e^{\beta(E_{\max}-E_0)} \right)^\alpha \right\} \end{aligned} \quad (80)$$

If the expression $\frac{u}{v} e^{\beta(E_{\max}-E_0)}$ is less than 1 then the maximum occurs at $\alpha = 0$, and otherwise the maximum occurs at $\alpha = n$. To determine this transition we set

$$\frac{u(\epsilon)}{v(\epsilon)} e^{\beta(E_{\max}-E_0)} = 1. \quad (81)$$

This can be rewritten as

$$\frac{3-\epsilon}{6-8\epsilon} e^{\beta(E_{\max}-E_0)} = 1. \quad (82)$$

The function $(3-\epsilon)/(6-8\epsilon)$ is increasing in ϵ , achieving its minimum value $1/2$ at $\epsilon = 0$ and its maximum value 1 at $\epsilon = 3/7$. We want to find in which cases there exists a threshold noise level ϵ_* at which the transition in Eq. (82) occurs.

For the case $E_{\max} = E_0$, namely when the negative component overlaps the highest energy state distribution, this threshold is constant in temperature and given by $\epsilon_* = 3/7$.

There is a threshold temperature value β_* given by

$$\beta_* := \frac{1}{E_{\max} - E_0} \ln 2. \quad (83)$$

Below the threshold temperature, $0 \leq \beta \leq \beta_*$, the transition is well defined and the threshold noise level at which it occurs is given by

$$\epsilon_*(\beta) := 3 - \frac{18}{8 - e^{(E_{\max}-E_0)\beta}}. \quad (84)$$

This encompasses the case $E_{\max} = E_0$.

We note that the threshold noise level ϵ_* becomes 0 when $\beta = \beta_*$, namely the difference between the maximum energy level and the level associated with the state negativity is equal to the energy extracted from a single particle Szilard's engine,

$$E_{\max} - E_0 = kT_* \ln 2. \quad (85)$$

For $\beta > \beta_*$, we have

$$\frac{3-\epsilon}{6-8\epsilon} e^{\beta(E_{\max}-E_0)} > \frac{3-\epsilon}{6-8\epsilon} 2 \geq \frac{1}{2} 2 = 1,$$

so there is no transition and the threshold noise level remains $\epsilon_* = 0$.

The maximum value of the Gibbs-rescaled distribution finally is

$$w(\rho_S|\gamma)_{\max} = \begin{cases} (3\mathcal{Z})^n v^n e^{n\beta E_0}, & \epsilon \leq \epsilon_*, \\ (3\mathcal{Z})^n u^n e^{n\beta E_{\max}}, & \epsilon > \epsilon_*. \end{cases} \quad (86)$$

Case (C1) corresponds to $(i, j, k) = (0, 0, 0)$, while case (C2) corresponds to the the conditions

$$(i, j, k) = \begin{cases} (0, n, 0), & \text{if } E_{\max} = E_1, \\ (0, 0, n), & \text{if } E_{\max} = E_2. \end{cases} \quad (87)$$

Note that if $E_{\max} = E_0$, we get $\epsilon_* = 3/7$ and (C1) holds for all ϵ and β . Assuming $E_{\max} > E_0$, we get the cases:

- $\beta = 0$. Then, $\epsilon_* = 3/7$ and (C1) holds for all noise levels ϵ . This is the unital fragment.
- $0 < \beta \leq \beta_*$. Then, $0 \leq \epsilon_* < 3/7$ and (C1) holds for $\epsilon \leq \epsilon_*$, while (C2) holds for $\epsilon \geq \epsilon_*$.

- $\beta > \beta_*$. Then, $\epsilon_* = 0$ and (C2) holds for all ϵ .

In (C1), the multiplicity is $m_{000} = 1$ and the corresponding Wigner components are $\mathbf{w}(\rho_S)_{000}, \mathbf{w}(\gamma)_{000}$. In (C2), we have $E_{\max} = E_1$ ($E_{\max} = E_2$), the multiplicity is $m_{0n0} = 3^n$ ($m_{00n} = 3^n$) and the corresponding Wigner components are $\mathbf{w}(\rho_S)_{0n0}, \mathbf{w}(\gamma)_{0n0}$ ($\mathbf{w}(\rho_S)_{00n}, \mathbf{w}(\gamma)_{00n}$). The first elbow coordinate expression can finally be expressed as

$$(x_0, L_0) = \begin{cases} \left(\left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^n, v^n \right), & \text{(C1)} \\ \left(\left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta} \right)^n, (3u)^n \right). & \text{(C2)} \end{cases} \quad (88)$$

Consider the distillation process in Eq. (67),

$$\rho_S(\epsilon)^{\otimes n} \xrightarrow{\mathcal{E} \in \mathcal{O}_{\gamma_\beta}} \rho_S(\epsilon')^{\otimes n'} \otimes \gamma_\beta^{\otimes (n-n')},$$

with $n > n'$ and $\epsilon > \epsilon'$. A necessary distillation constraint can be defined at the location x_* of the initial state's first elbow, expressed as

$$L_{\rho_S(\epsilon)^{\otimes n} | \gamma_\beta}(x_*) \geq L_{\rho_S(\epsilon')^{\otimes n'} | \gamma_\beta}(x_*). \quad (89)$$

The Lorenz curves of the initial and target states may each be described by either (C1) or (C2), depending on the physical parameters $\epsilon, \epsilon', \beta$. Specifically, we can get three scenarios:

1. (C1) \rightarrow (C1) if $E_{\max} = E_0$ OR $E_{\max} > E_0, \beta < \beta_*$ and $\epsilon' < \epsilon \leq \epsilon_*$.
2. (C2) \rightarrow (C1) if $E_{\max} > E_0, \beta < \beta_*$ and $\epsilon' \leq \epsilon_* < \epsilon$.
3. (C2) \rightarrow (C2) if $E_{\max} > E_0, \beta \geq \beta_*$.

Note that (C1) \rightarrow (C2) is impossible because it would imply $\beta < \beta_*$ and $\epsilon \leq \epsilon_* \leq \epsilon' < \epsilon$, a contradiction.

In all three scenarios, it is simple to check that the first elbow of the initial state is always located closer to 0 compared to the first elbow of the target state, $x_0 \leq x'_0$:

1. (C1) \rightarrow (C1). It is always true that $e^{-\beta E_0}/\mathcal{Z}_\beta \leq 1$, so

$$\begin{aligned} x_0 &= \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^n = \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n-n'} \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n'} \\ &< \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n'} = x'_0 \end{aligned}$$

2. (C2) \rightarrow (C1). It is always true that

$$\begin{aligned} \mathcal{Z}_\beta^{n-n'} &> (e^{-\beta E_{\max}})^{n-n'} = e^{-\beta(n-n')E_{\max}}, \text{ so} \\ x_0 &= \left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta} \right)^n = \frac{e^{-\beta(nE_{\max}-n'E_0)}}{\mathcal{Z}_\beta^{n-n'}} \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n'} \\ &< \frac{e^{-\beta(nE_{\max}-n'E_0)}}{e^{-\beta(n-n')E_{\max}}} \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n'} \\ &= e^{-\beta n'(E_{\max}-E_0)} \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n'} \\ &\leq \left(\frac{e^{-\beta E_0}}{3\mathcal{Z}_\beta} \right)^{n'} = x'_0. \end{aligned}$$

3. (C2) \rightarrow (C2). It is always true that $e^{-\beta E_{\max}}/\mathcal{Z}_\beta \leq 1$, so

$$\begin{aligned} x_0 &= \left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta} \right)^n = \left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta} \right)^{n-n'} \left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta} \right)^{n'} \\ &\leq \left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta} \right)^{n'} = x'_0 \end{aligned}$$

We can thus find the target state Lorenz curve coordinate L'_* at location x_* by interpolating between the origin and the target state's first elbow,

$$L'_* = \frac{x_*}{x'_0} L'_0. \quad (90)$$

We need $L_0 \geq L'_*$. Therefore, restricting to the first elbow, the necessary distillation condition in Eq. (89) simply becomes

$$\frac{L_0}{x_0} \geq \frac{L'_0}{x'_0}. \quad (91)$$

Involving more elbows gives stricter, but more convoluted necessary distillation constraints.

We now compute distillation bounds based on the first elbow constraint in the three possible scenarios. [\[Does anything special happen if we have degeneracies in energy?\]](#) [\[If we have degeneracies, basically the Gibbs-rescaled distribution in Eq. \(73\) has some equal components, so Lorenz curves have fewer elbows, but because the first elbow constraint boils down to a ratio comparison Eq. \(91\) it doesn't change the analysis\]](#) We express them in terms of the free energy of thermal state γ_β denoted by

$$F_\beta := \text{tr}[\gamma_\beta H] + \beta^{-1} \text{tr}[\gamma_\beta \ln \gamma_\beta] = -\beta^{-1} \ln \mathcal{Z}_\beta. \quad (92)$$

Substituting coordinates from Eq. (88) appropriately in Eq. (91), we get the following necessary conditions:

1. (C1) \rightarrow (C1) if $E_{\max} = E_0$ OR $E_{\max} > E_0, \beta < \beta_*$ and $\epsilon' < \epsilon \leq \epsilon_*$.

$$\frac{n}{n'} \geq \frac{\ln(1 - \frac{4}{3}\epsilon') + \beta(E_0 - F_\beta)}{\ln(1 - \frac{4}{3}\epsilon) + \beta(E_0 - F_\beta)}, \text{ or} \quad (93)$$

$$\epsilon \leq \frac{3}{4} - \frac{3}{4} \left(1 - \frac{4}{3}\epsilon'\right)^{\frac{n'}{n}} \left(\frac{e^{-\beta E_0}}{\mathcal{Z}_\beta}\right)^{1 - \frac{n'}{n}} \quad (94)$$

This leads to an upper bound for the distillation rate $R = n'/n$ of pure Strange states ($\epsilon' = 0$),

$$R \leq R(\epsilon, \beta) := 1 - \frac{\ln\left(\frac{1}{1 - \frac{4}{3}\epsilon}\right)}{\beta(E_0 - F_\beta)}, \quad (95)$$

2. (C2) \rightarrow (C1) if $E_{\max} > E_0$, $\beta < \beta_\star$ and $\epsilon' \leq \epsilon_\star < \epsilon$.

$$\frac{n}{n'} \geq \frac{\ln(1 - \frac{4}{3}\epsilon') + \beta(E_0 - F_\beta)}{\ln(\frac{1}{2} - \frac{1}{6}\epsilon) + \beta(E_{\max} - F_\beta)}, \text{ or} \quad (96)$$

$$\epsilon \leq 3 - 6 \left(1 - \frac{4}{3}\epsilon'\right)^{\frac{n'}{n}} \frac{e^{-\beta E_{\max}} e^{\frac{n'}{n}\beta E_0}}{\mathcal{Z}_\beta^{1 - \frac{n'}{n}}} \quad (97)$$

This leads to an upper bound for the distillation rate $R = n'/n$ of pure Strange states ($\epsilon' = 0$),

$$R \leq R(\beta) := 1 - \frac{\ln\left(\frac{1}{1 - \frac{1}{3}\epsilon}\right) + (E_{\max} - E_0)(\beta_\star - \beta)}{\beta(E_0 - F_\beta)}, \quad (98)$$

3. (C2) \rightarrow (C2) if $E_{\max} > E_0$, $\beta \geq \beta_\star$.

$$\frac{n}{n'} \geq \frac{\ln(\frac{1}{2} - \frac{1}{6}\epsilon') + \beta(E_{\max} - F_\beta)}{\ln(\frac{1}{2} - \frac{1}{6}\epsilon) + \beta(E_{\max} - F_\beta)}, \text{ or} \quad (99)$$

$$\epsilon \leq 3 - 6 \left(\frac{1}{2} - \frac{1}{6}\epsilon'\right)^{\frac{n'}{n}} \left(\frac{e^{-\beta E_{\max}}}{\mathcal{Z}_\beta}\right)^{1 - \frac{n'}{n}} \quad (100)$$

This leads to an upper bound for the distillation rate $R = n'/n$ of pure Strange states ($\epsilon' = 0$),

$$R \leq R(\beta) := 1 - \frac{\ln\left(\frac{1}{1 - \frac{1}{3}\epsilon}\right)}{-\ln 2 + \beta(E_{\max} - F_\beta)}, \quad (101)$$

Consider pure Strange state distillation at high temperatures, $\beta \leq \beta_\star$. Depending on the noise level of the initial state, we have either a (C1) \rightarrow (C1) or a (C2) \rightarrow (C1) scenario.

At every $\beta > 0$, the rates indicate a transition at ϵ_\star . As β increases, the rate becomes stricter. At the high temperature limit, $\beta \rightarrow 0$, we get $\beta F_\beta \rightarrow -\ln 3$, and we retrieve the unital fragment bound,

$$R(\beta) \xrightarrow{\beta \rightarrow 0} \frac{\ln(3 - 4\epsilon)}{\ln 3}. \quad (102)$$

Consider pure Strange state distillation at low temperatures, $\beta \rightarrow \infty$. Independently of the noise level of the initial state, we have a (C2) \rightarrow (C2) scenario.

As β increases, the rate becomes looser. At the zero temperature limit, F_β tends to zero, so $R(\epsilon, \beta) \xrightarrow{\beta \rightarrow \infty} 1$. \square

VI. COMPARISON WITH EXISTING DISTILLATION RATES

We have derived a *distillation bound* $R(\epsilon, \epsilon', \beta)$ for the process

$$\rho_S(\epsilon)^{\otimes n} \longrightarrow \rho_S(\epsilon')^{\otimes n'} \otimes \gamma_\beta^{\otimes (n-n')}, \quad (103)$$

with even n, n' and $0 \leq \epsilon' < \epsilon \leq 3/7$, where the n -copy, ϵ -noisy Strange state is

$$\rho_S(\epsilon)^{\otimes n} := \left[(1-\epsilon) |S\rangle \langle S| + \epsilon \frac{1}{3} \mathbb{1} \right]^{\otimes n}. \quad (104)$$

The bound informs us that any distillation protocol in the γ_β -fragment with an ϵ -noisy input and ϵ' -noisy output has a *distillation rate*

$$\frac{n'}{n} \leq R(\epsilon, \epsilon', \beta). \quad (105)$$

We apply our results on the ternary Golay code distillation protocol suggested in [31]. In this protocol, a single Strange state is distilled with noise level

$$\epsilon'(n, \epsilon) \approx \frac{1}{\alpha} (\alpha \epsilon)^{n^\xi}, \quad (106)$$

where $\alpha \approx 1.75$ and $\xi \approx 0.112$ are constants specific to the protocol. This can be rewritten in terms of the protocol distillation rate,

$$R_{\text{Golay}}(\epsilon, \epsilon') = \frac{1}{n} \approx \left(\frac{\log(\alpha \epsilon)}{\log(\alpha \epsilon')} \right)^{\frac{1}{\xi}} \quad (107)$$

We compare the distillation rate of the protocol with our distillation bound,

$$R_{\text{Golay}}(\epsilon, \epsilon') \leq R(\epsilon, \epsilon', \beta) \quad (108)$$

in figure Fig. (8)

We can also rewrite the expression as

$$\epsilon_{\text{Golay}}(\epsilon', n) \approx \frac{1}{\alpha} (\alpha \epsilon')^{n^{-\xi}}, \quad (109)$$

and compare with the numerically optimal error bounds provided by majorisation as shown in Fig. (9)

A recent distillation bound was introduced in [33], where they show that for the deterministic conversion of n copies of a qubit state ρ to a target pure magic state $|\psi\rangle$, the following bound on the distillation rate holds,

$$\frac{n'}{n} \leq \frac{\log \Lambda^+(\rho)}{\log F(\psi)^{-1}}, \quad (110)$$

where $\Lambda^+(\rho)$ is the generalised robustness of state ρ , calculable via a costly SDP and $F(\psi)$ denotes the stabiliser fidelity of state $|\psi\rangle$. They demonstrate the performance of their bound by considering the purification process of the qubit magic state $|H\rangle$, with their figure replicated

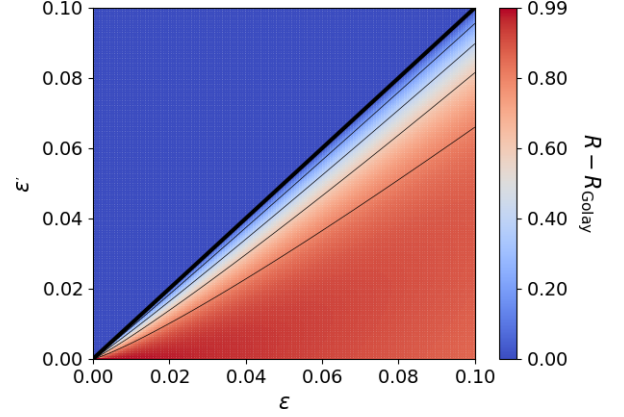


FIG. 8. **Comparison with ternary Golay distillation rate.** Our bound is not violated but gets close to saturated for some noise levels. The thick diagonal line indicates where $\epsilon' = \epsilon$.

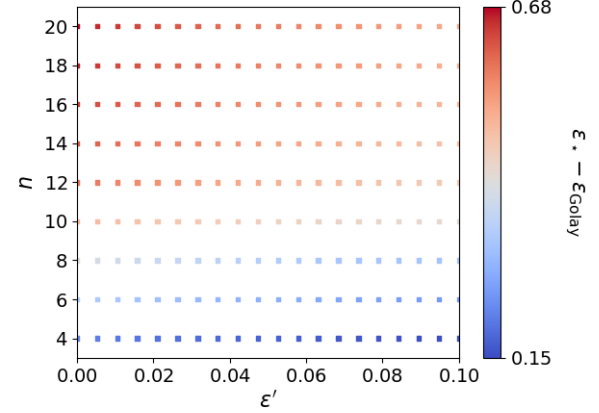


FIG. 9. **Comparison with ternary Golay output error rate.** Our bound is not violated but achieves a maximum saturation of $\epsilon_* - \epsilon_{\text{Golay}} = 0.147$.

in Fig. (10). For high initial noise level ($\epsilon = 0.25$), their bound dictates that $R = 1/150$, which is pretty good.

Similar concrete rates have been studied for qudits in two occasions: First, the mana bound, which we have demonstrated in Fig. (5). Secondly, Wang *et al.*'s max-thauma bound [35], which as far as I can tell, is less tight than mana (?). It is defined via an SDP, but we can use its properties (super-additivity and zero at free states) to readily get a bound on noisy Strange states,

$$\theta_{\max}(\rho_S(\epsilon)) \geq (1-\epsilon) |S\rangle \langle S| + \epsilon \frac{1}{3} \mathbb{1} = (1-\epsilon) \theta_{\max}(\rho_S(0)). \quad (111)$$

Their bound simply states that any deterministic single-copy distillation process $\rho^{\otimes n} \longrightarrow \tau$ requires an initial number of copies $n \geq \theta_{\max}(\tau)/\theta_{\max}(\rho)$. Therefore, the rate of our pure Strange state distillation process Eq. (36)

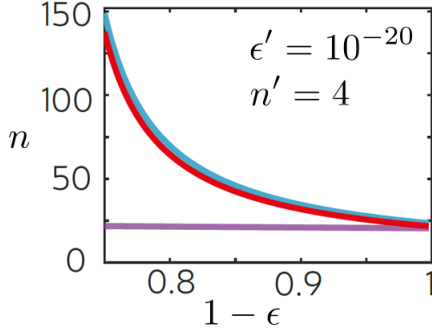


FIG. 10. **Comparison of Seddon *et al.*'s bound [33] in Eq. (110) with a generic resource bound in [34].** The process they consider is purifying n copies of $(1 - \epsilon) |H\rangle\langle H| + \epsilon/3 \mathbb{1}$ to n' copies of $(1 - \epsilon') |H\rangle\langle H| + \epsilon'/3 \mathbb{1}$. Their bound (red and blue) is pretty good compared with the generic (violet) bound. Not sure why it looks like their bound gives $n \rightarrow 25$ as $1 - \epsilon \rightarrow 1$.

can only get as tight as

$$R = 1/n = \frac{\theta_{\max}(\rho_S(\epsilon))}{\theta_{\max}(\rho_S(0))} \leq 1 - \epsilon, \quad (112)$$

which is less tight than mana as could be readily observed if plotted in Fig. (5).

[Going forward:

1. Does the “last elbow” give the same bound? Is it as simple?
2. Numerically is this about as good as it gets from majorisation?
3. How does this rate compare with existing literature?

VII. CONCLUSION

1. Introduced fragments
2. Identify symmetries of the setup
3. Combined single-shot thermodynamics with magic
4. Can we solve other cases exactly? (apart from single qutrit)

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Appendix A: Properties of Wigner distributions

Here we present important properties of the Wigner distribution that are used throughout the paper.

Proposition 11. *The Wigner distribution of a state $\rho \in \mathcal{B}(\mathcal{H}_d)$ is*

- (i) *Real valued:* $W_\rho \in \mathbb{R}^{d^2}$;
- (ii) *Normalised:* $\sum_{\mathbf{z} \in \mathcal{P}_d} W_\rho(\mathbf{z}) = 1$;
- (iii) *Bounded:* $|W_\rho(\mathbf{x})| \leq \frac{1}{d}$.
- (iv) *Additive under mixing:*
 $W_{\sum_i p_i \rho_i}(\mathbf{x}) = \sum_i p_i W_{\rho_i}(\mathbf{x})$;

(v) *Multiplicative under tensor products:*

$$W_{\rho_A \otimes \rho_B}(\mathbf{x}_A \oplus \mathbf{x}_B) = W_{\rho_A}(\mathbf{x}_A) W_{\rho_B}(\mathbf{x}_B).$$

Not all distributions satisfying these properties correspond to positive semi-definite states.

Proof. [Proof?] □

Proposition 12. *The Wigner distribution of a CPTP operation $\mathcal{E} : \mathcal{B}(\mathcal{H}_{d_A}) \mapsto \mathcal{B}(\mathcal{H}_{d_B})$ is*

- 1. *Real-valued:* $W_{\mathcal{E}} \in \mathbb{R}^{d^2} \times \mathbb{R}^{d^2}$;
- 2. *Normalised:* $\sum_{\mathbf{z} \in \mathcal{P}_{d_B}} W_{\mathcal{E}}(\mathbf{z}|\mathbf{x}) = 1$
for any $\mathbf{x} \in \mathcal{P}_{d_A}$;
- 3. *Bounded:* $|W_{\mathcal{E}}(\mathbf{y}|\mathbf{x})| \leq \frac{d_A}{d_B}$;
- 4. *Transitive:* $W_{\mathcal{E}(\rho)}(\mathbf{y}) = \sum_{\mathbf{z} \in \mathcal{P}_{d_A}} W_{\mathcal{E}}(\mathbf{y}|\mathbf{z}) W_\rho(\mathbf{z})$
for any $\mathbf{y} \in \mathcal{P}_{d_B}$.

If $d_A = d_B$, and in particular if operation \mathcal{E} maps a Hilbert space onto itself, then the stochasticity condition $|W_{\mathcal{E}}(\mathbf{y}|\mathbf{x})| \leq 1$ is satisfied.

Proof. [Proof?] □

A $(d \times d)$ matrix S is in $S_d(\mathbf{d})$, if

- 1. $S_{ij} \geq 0$ for all $i, j \in \mathbb{Z}_d$;
- 2. $\sum_{j=1}^n S_{ij} = 1$ for all $i \in \mathbb{Z}_d$;
- 3. $S\mathbf{d} = \mathbf{d}$.

Appendix B: Properties of majorization

1. Simple-majorization equivalence conditions

In the unital fragment, namely the limit of infinite temperature, $\beta = 0$, the free state is the maximally mixed state $\frac{1}{d}\mathbb{1}$ with uniform Wigner distribution $\frac{1}{d}\mathbf{1} = (\frac{1}{d}, \dots, \frac{1}{d})$. This fragment is governed by simple majorization and we first prove strong equivalences for this type of majorization.

Theorem 13. *Given $\mathbf{x}, \mathbf{y} \in \mathbb{R}^n$ the following statements are equivalent:*

- 1. $\mathbf{x} \prec \mathbf{y}$;
- 2. $\mathbf{x} = B\mathbf{y}$ for bistochastic B ;
- 3. $\sum_{i=1}^n |x_i - t| \leq \sum_{i=1}^n |y_i - t|$ for all $t \in \mathbb{R}$;
- 4. $\sum_{i=1}^n (x_i - t)^+ \leq \sum_{i=1}^n (y_i - t)^+$ for all $t \in \mathbb{R}$ and $\sum_{i=1}^n x_i = \sum_{i=1}^n y_i$, where $(x)^+ = \max\{x, 0\}$;
- 5. $L_{\mathbf{x}}(k) \leq L_{\mathbf{y}}(k)$ for $k = 1, \dots, n-1$ and $L_{\mathbf{x}}(n) = L_{\mathbf{y}}(n)$.

Proof. [Proof?] □

2. Embedding map

Any \mathbf{d} -majorization problem can be rephrased as a simple majorization problem in a higher dimensional space via the embedding map.

Definition 14. *The embedding map $\Gamma_{\mathbf{d}} : \mathbb{R}^n \mapsto \mathbb{R}^D$, $D = \sum_{i=1}^n d_i$ is the function*

$$\Gamma_{\mathbf{d}}(\mathbf{w}) = \bigoplus_{i=1}^n w_i \frac{1}{d_i} \mathbf{1}, \quad (\text{B1})$$

where $\mathbf{1}/d_i$ is the d_i -dimensional uniform distribution. The left inverse $\Gamma_{\mathbf{d}}^{-1} : \mathbb{R}^D \mapsto \mathbb{R}^n$ is defined to sum up the elements in each block of $\Gamma_{\mathbf{d}}(\mathbf{w})$, so that

$$\Gamma_{\mathbf{d}}^{-1}(\oplus_{i=1}^n w_i \mathbf{1}/d_i) = \mathbf{w}. \quad (\text{B2})$$

This is not a right inverse, because $\Gamma_{\mathbf{d}}$ is not surjective.

The direct sum simply amounts to listing the uniform distributions one after the other. The embedding map maps the Gibbs distribution to the uniform distribution, $\Gamma_{\mathbf{d}}(\mathbf{d}) = \mathbf{1}/D$. Then, a non-increasing ordering $\Gamma_{\mathbf{d}}(\mathbf{z})^\downarrow$ in the new space, corresponds to the so-called “ β -ordering” of the original vector denoted by the permutation π in Definition 3, mapping $(w_i/d_i) \mapsto (w_i/d_i)^\downarrow$ for all $i = 1, \dots, n$.

3. d -majorization equivalence conditions

Theorem 15. *Given $\mathbf{x}, \mathbf{y}, \mathbf{d} \in \mathbb{R}^n$, such that the components of \mathbf{d} are positive, the following statements are equivalent:*

1. $\mathbf{x} \prec_{\mathbf{d}} \mathbf{y}$;
2. $\Gamma_{\mathbf{d}}(\mathbf{x}) \prec \Gamma_{\mathbf{d}}(\mathbf{y})$;
3. $\sum_{i=1}^n |x_i - td_i| \leq \sum_{i=1}^n |y_i - td_i|$ for all $t \in \mathbb{R}$;
4. $\sum_{i=1}^n (x_i - td_i)^+ \leq \sum_{i=1}^n (y_i - td_i)^+$ for all $t \in \mathbb{R}$ and $\sum_{i=1}^n x_i = \sum_{i=1}^n y_i$;
5. $L_{\mathbf{x}|\mathbf{d}}(k) \leq L_{\mathbf{y}|\mathbf{d}}(k)$ for $k = 1, \dots, n-1$ and $L_{\mathbf{x}|\mathbf{d}}(n) = L_{\mathbf{y}|\mathbf{d}}(n)$.

Proof. [Proof?] □

Appendix C: Technical properties of σ -fragments

[this section hasn't been refereed in various parts of the main text]

Theorem 16. *Let $\mathcal{R}_\sigma = (\mathcal{O}_\sigma, \mathcal{F}_\sigma)$ be a σ -fragment of magic theory $\mathcal{R} = (\mathcal{O}, \mathcal{F})$. The following statements hold:*

1. No σ -fragment is empty.
2. If a free operation leaves two states invariant, then it also leaves their mixtures invariant,

$$\mathcal{O}_\sigma \cap \mathcal{O}_{\sigma'} \subseteq \mathcal{O}_{p\sigma + (1-p)\sigma'} \text{ for all } p \in [0, 1].$$

3. Let \mathcal{E} be a CPTP operation with Wigner distribution $W_\mathcal{E}$. For $\mathcal{R} = \mathcal{R}_{\max}$, $\mathcal{E} \in \mathcal{O}_\sigma$ iff $W_\mathcal{E} \in S_{d^2}(W_\sigma)$. [I don't like this one]

Proof.

1. The identity channel $1_{\mathcal{D}} : \mathcal{D} \mapsto \mathcal{D}$ belongs to every σ -fragment, as $1_{\mathcal{D}} \in \mathcal{O}$ and $1_{\mathcal{D}}\sigma = \sigma$ for all $\sigma \in \mathcal{F}$.
2. Let $\mathcal{E} \in \mathcal{O}_\sigma \cap \mathcal{O}_{\sigma'}$. Then $\mathcal{E} \in \text{CPTP}$ and corresponds to stochastic Wigner distribution $W_\mathcal{E}$ such that $W_\mathcal{E}W_\sigma = W_\sigma$ and $W_\mathcal{E}W_{\sigma'} = W_{\sigma'}$. Then, $W_\mathcal{E}W_{p\sigma + (1-p)\sigma'} = W_{p\sigma + (1-p)\sigma'}$ for any $p \in [0, 1]$ due to the additive property 11 of the Wigner distribution, implying that state $p\sigma + (1-p)\sigma'$ is also left invariant by \mathcal{E} .

3. Let $\mathcal{O}'_\sigma := \{\mathcal{E} \in \text{CPTP} : W_\mathcal{E} \in S_{d^2}(W_\sigma)\}$ be the described set of operations.

Suppose \mathcal{E} is in \mathcal{O}_σ , then $\mathcal{E} \in \text{CPTP}$ and $W_\mathcal{E} \in S_{d^2}(W_\sigma)$ due to property 4 of Proposition 12, hence $\mathcal{O}_\sigma \subseteq \mathcal{O}'_\sigma$.

Conversely, suppose $\mathcal{E} \in \text{CPTP}$ with $W_\mathcal{E} \in S_{d^2}(W_\sigma)$. Then, $W_\mathcal{E}(\mathbf{y}|\mathbf{x}) \geq 0$ for all \mathbf{x}, \mathbf{y} , hence $\mathcal{E} \in \mathcal{O}$. Furthermore, $W_\mathcal{E}W_\sigma = W_\sigma$ implies $\mathcal{E}(\sigma) = \sigma$ using Eq. (12) defined for any CPTP operation \mathcal{E} . Hence, $\mathcal{O}'_\sigma \subseteq \mathcal{O}_\sigma$. □

Any free state $\sigma \in \mathcal{F}$ corresponds to a d^2 -dimensional probability distribution W_σ and any free operation $\mathcal{E} \in \mathcal{O}$ corresponds to a $d^2 \times d^2$ stochastic matrix (or conditional probability distribution) $W_\mathcal{E}$. Note that these mappings are one-to-one due to the orthogonality of the phase-point operators as an operator basis.

Note further that free states \mathcal{F} are mapped onto a strict subset of the set of probability distributions. As a counterexample, the sharp d^2 -dimensional probability distribution $(1, 0, \dots, 0)$ does not correspond to any qudit Wigner distribution because of the boundedness condition in Proposition 11.

Similarly, not all stochastic matrices correspond to completely positive operations. As an example, consider the permutation matrix

$$\Pi_X = \begin{pmatrix} 0 & 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 \\ 0 & 0 & 0 & 1 & 0 \\ 1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 & 0 \end{pmatrix} \otimes \begin{pmatrix} 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 \\ 0 & 0 & 0 & 1 & 0 \\ 1 & 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 \end{pmatrix} \in S_5(W_{\frac{1}{5}\mathbb{1}}). \quad (\text{C1})$$

It preserves the uniform distribution $W_{\frac{1}{5}\mathbb{1}}$, but it does not correspond to any positive (hence quantum) operation.

Appendix D: Properties of Lorenz curves

[Prove statements underneath Definition 3?]

1. Basic properties

Furthermore, single-copy Lorenz curves are simply additive in noise,

$$L_{(1-\epsilon)\rho + \epsilon\sigma}(x) = (1-\epsilon)L_{\rho| \sigma}(x) + \epsilon L_{\sigma| \sigma}(x), \quad (\text{D1})$$

for any σ, ρ . This is not true for higher number of copies.

2. σ -fragment continuity

Proposition 17. *The Lorenz curve $L_{\rho| \sigma}(x)$ of a state ρ with respect to state σ is uniformly continuous in σ .*

Proof. [Idea:] If σ and σ' are δ -close with respect to a state norm, then the curves $L_{\rho| \sigma}(x)$ and $L_{\rho| \sigma'}(x)$ are ϵ -close at all $x \in [0, 1]$. □

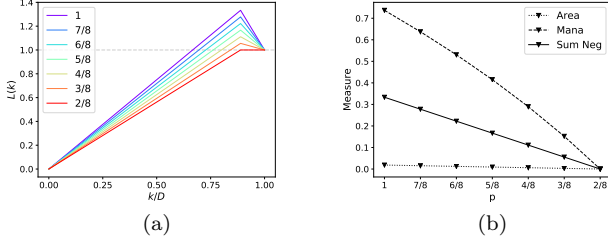


FIG. 11. (a) Lorenz curve of $p|S\rangle\langle S| + (1-p)\frac{1}{d}\mathbb{1}$ for p given in the legend. (b) Different measures for the states on the left. [can replace Fig. (3) with a more informative version of this figure]

3. Area monotone

[I don't believe in the area monotone - should we keep?]

Let $L_{>1}$ be the set of points on the Lorenz curve $L_{\rho|\sigma}(k)$ that lie above 1 of state ρ in the σ -fragment. If ρ is a free state, $L_{>1}$ is empty and $\mathcal{A}_\sigma(\rho) = 0$. Otherwise, $\mathcal{A}_\sigma(\rho) > 0$ and it can be calculated exactly using the trapezium rule or the shoelace formula. Let k be the index of the first point $(x_k, L_{\rho|\sigma}(k))$ lying above 1. Then $L_{\rho|\sigma}(k)$ crosses 1 at

$$x_{\text{int}} = x_k - \frac{x_k - x_{k-1}}{L_{\rho|\sigma}(k) - L_{\rho|\sigma}(k-1)} L_{\rho|\sigma}(k), \quad (\text{D2})$$

as well as at $(x_d, L_{\rho|\sigma}(d)) = (1, 1)$

Now we can define

$$L_{>1}^+ = \{(x_i, y_i)\}_{0 \leq i \leq n} \quad (\text{D3})$$

such that it contains the initial point of intersection $(x_0, y_0) \equiv (x_{n+1}, y_{n+1}) := (x_{\text{int}}, 1)$, all points (x_i, y_i) , labelled by $i = 1, \dots, n-1$ that lie above $y = 1$ contained in $L_{>1}$ and finally the second point of intersection $(x_n, y_n) = (1, 1)$. Then,

$$\mathcal{A}_\sigma(\rho) = \frac{1}{2} \sum_{i=0}^n (x_{i+1} y_i - x_i y_{i+1}). \quad (\text{D4})$$

4. Lorenz curve independent constraints

Theorem 18. Let ρ, τ be two quantum states with Lorenz curves $L_{\rho|\sigma}(x), L_{\tau|\sigma}(x)$ in the σ -fragment.

Let t be the number of elbows of $L_{\tau|\sigma}(x)$ at locations x_1, \dots, x_t .

Then, $L_{\rho|\sigma}(x) \geq L_{\tau|\sigma}(x)$ for all $x \in [0, 1]$ iff $L_{\rho|\sigma}(x_i) \geq L_{\tau|\sigma}(x_i)$ for all $i = 1, \dots, t$.

Proof. $L_{\rho|\sigma}(x) \geq L_{\tau|\sigma}(x)$ for all $x \in [0, 1]$ trivially implies $L_{\rho|\sigma}(x_i) \geq L_{\tau|\sigma}(x_i)$ for all $i = 1, \dots, n'$.

Conversely, assume that $L_{\rho|\sigma}(x_i) \geq L_{\tau|\sigma}(x_i)$ for all $i = 1, \dots, r$. First, let $x_0 = 0$ and $x_{n'+1} = 1$,

so that $L_{\rho|\sigma}(x_0) = L_{\tau|\sigma}(x_0) = 0$ and $L_{\rho|\sigma}(x_{n'+1}) = L_{\tau|\sigma}(x_{n'+1}) = 1$. Hence, we can extend the set of elbows E to $E' = E \cup \{x_0, x_{n'+1}\}$.

Pick two consecutive locations x_i, x_{i+1} in E' and consider the line segment $\ell_\tau(x)$ connecting points $(x_i, L_{\tau|\sigma}(x_i))$ and $(x_{i+1}, L_{\tau|\sigma}(x_{i+1}))$ as well as the line segment $\ell_\rho(x)$ connecting points $(x_i, L_{\rho|\sigma}(x_i))$ and $(x_{i+1}, L_{\rho|\sigma}(x_{i+1}))$. This is illustrated in Fig. (12).

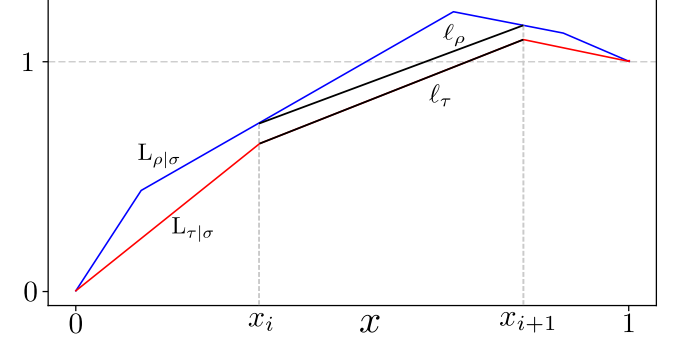


FIG. 12. Illustration of Theorem 18.

Due to concavity of $L_{\rho|\sigma}$, it is clear that for all $x \in [x_i, x_{i+1}]$, we have $L_{\rho|\sigma}(x) \geq \ell_\rho(x) \geq \ell_\tau(x) = L_{\tau|\sigma}(x)$. This argument can be made in all intervals $[x_i, x_{i+1}]$ with $i = 0, \dots, n'$, so the proof is complete. \square

This theorem is of practical importance in calculating the necessary distillation constraints derived via majorization.

[In particular], it can be used in the discussion of standard Lorenz curves to set the necessary constraints,

$$\begin{aligned} L(x_{I'(0, m'_0, 9^{n-n'})}) &\geq L(x_{I'(0, m'_0, 9^{n-n'})}), \\ L(x_{I'(1, m'_1, 9^{n-n'})}) &\geq L(x_{I'(1, m'_1, 9^{n-n'})}), \\ &\vdots \\ L(x_{I'(n', m'_{n'}, 9^{n-n'})}) &\geq L(x_{I'(n', m'_{n'}, 9^{n-n'})}), \end{aligned}$$

where $x_{I'(0, m'_0, 9^{n-n'})}, x_{I'(1, m'_1, 9^{n-n'})}, \dots, x_{I'(n', m'_{n'}, 9^{n-n'})} \in T$ are the locations of the elbows of $L_{\rho_S(\epsilon') \otimes n' | \frac{1}{d} \mathbb{1}}$ (excluding 0 and 1).

Appendix E: Technical derivation for Lorenz curves in unital fragment

1. Binomial distributions and error bounds

Consider an experiment consisting of n trials of throwing a p -coin, that is a coin with probability p of landing on one side and $1 - p$ of landing on the other. This scenario describes the construction of an n -copy Strange state Wigner distribution from a single-copy as presented in ???. The experiment is described by binomial statistics

and we write down the sum over an even number of successful trials Φ_+ (over an odd number of successful trials Φ_-),

$$\Phi_+(af; nf, p) := \sum_{\ell=0}^{af/2} \binom{nf}{2\ell} p^{2\ell} (1-p)^{nf-2\ell},$$

for even integers $af \in [0, nf]$, (E1)

$$\Phi_-(af; nf, p) := \sum_{\ell=1}^{(af-1)/2} \binom{nf}{2\ell+1} p^{2\ell+1} (1-p)^{nf-(2\ell+1)},$$

for odd integers $af \in [0, nf]$. (E2)

Note that index af only takes even (odd) values when labelling Φ_+ (Φ_-). In Appendix E 2, we will use Φ_+ and Φ_- to express the elbow coordinates of standard Lorenz curves.

We define the classical entropy of a p -coin as well as the classical relative entropy between a p -coin and a q -coin,

$$S(p) := -p \log p - (1-p) \log(1-p), \quad (\text{E3})$$

$$S(p||q) := p \log \frac{p}{q} + (1-p) \log \frac{1-p}{1-q}. \quad (\text{E4})$$

They are symmetric in the sense that $S(p) = S(1-p)$ and $S(p||q) = S(1-p||1-q)$.

A useful result is the entropic bound on a combination [36].

Lemma 19. For all $\ell \in [1, np-1]$,

$$\left[8\ell \left(1 - \frac{\ell}{np} \right) \right]^{-\frac{1}{2}} 2^{nS(\frac{\ell}{np})} \leq \binom{np}{\ell} \leq \quad (\text{E5})$$

$$\left[2\pi\ell \left(1 - \frac{\ell}{np} \right) \right]^{-\frac{1}{2}} 2^{nS(\frac{\ell}{np})}. \quad (\text{E6})$$

Proof. For $\ell = 1, 2, np-1, np-2$ check by direct calculation. For all other cases, use Stirling's approximation. [CITE] \square

With the help of this lemma, we directly arrive at

Theorem 20. Given fixed $n > 0$ and p , Φ_+, Φ_- satisfy

the following bounds:

1. $\Phi_+(a; n, p) \geq \sum_{\ell=0}^{np/2} \left[16\ell \left(1 - \frac{2\ell}{np} \right) \right]^{-\frac{1}{2}} 2^{-nS(\frac{2\ell}{np}||p)},$
for all even $a \in [2, np]$
2. $\Phi_+(a; n, p) \leq \sum_{\ell=0}^{np/2} \left[4\pi\ell \left(1 - \frac{2\ell}{np} \right) \right]^{-\frac{1}{2}} 2^{-nS(\frac{2\ell}{np}||p)},$
for all even $a \in [2, np]$
3. $\Phi_-(a; n, p) \geq \sum_{\ell=1}^{(np-1)/2} \left[16(\ell+1) \left(1 - \frac{2\ell+1}{np} \right) \right]^{-\frac{1}{2}} \times$
 $\times 2^{-nS(\frac{2\ell+1}{np}||p)},$ for all odd $a \in [1, np]$
4. $\Phi_-(a; n, p) \leq \sum_{\ell=1}^{(np-1)/2} \left[4\pi(\ell+1) \left(1 - \frac{2\ell+1}{np} \right) \right]^{-\frac{1}{2}} \times$
 $\times 2^{-nS(\frac{2\ell+1}{np}||p)},$ for all odd $a \in [1, np]$

Proof. All four statements follow from application of Lemma 19 on the combinatorial coefficient and the definition of relative entropy given in Eq. (E4) \square

2. Standard Lorenz curve elbow coordinates

The Wigner distribution of the n -copy qutrit maximally mixed state $(\mathbb{1}/3)^{\otimes n}$ is the uniform distribution

$$W_{(\mathbb{1}/3)^{\otimes n}} = \left(\overbrace{\frac{1}{9^n}, \dots, \frac{1}{9^n}}^{9^n} \right). \quad (\text{E7})$$

The Wigner distribution of the 1-copy ϵ -noisy Strange state $\rho_S(\epsilon)$ in the unital fragment is a permutation of

$$W_{\rho_S(\epsilon)} = \left(\overbrace{\frac{1}{6} - \frac{1}{18}\epsilon, \dots, \frac{1}{6} - \frac{1}{18}\epsilon}^8, \overbrace{-\frac{1}{3} + \frac{4}{9}\epsilon}^1 \right) \quad (\text{E8})$$

The two distinct components are plotted Fig. (13) as a function of noise. It is clear that in the unital fragment the Strange state contains Wigner negativities in the regime $0 \leq \epsilon \leq \frac{3}{4}$.

The Wigner distribution of the n -copy ϵ -noisy Strange state $\rho_S(\epsilon)^{\otimes n}$ in the unital fragment is given by the convolution $W_{\rho_S(\epsilon)^{\otimes n}} = W_{\rho_S(\epsilon)}^{\otimes n}$. In general, $\rho_S(\epsilon)^{\otimes n}$ contains $n+1$ distinct components, labelled $0, \dots, n$. We present the distinct Wigner components of $\rho_S(\epsilon)^{\otimes n}$ along with their multiplicities in Table I. Note that LHS (RHS) refers to elbow coordinates i on the left of and including (right of) the maximum, precisely

$$\text{LHS: } 0 \leq i \leq \left\lfloor \frac{n}{2} \right\rfloor; \quad (\text{E9})$$

$$\text{RHS: } \left\lfloor \frac{n}{2} \right\rfloor + 1 \leq i \leq n. \quad (\text{E10})$$

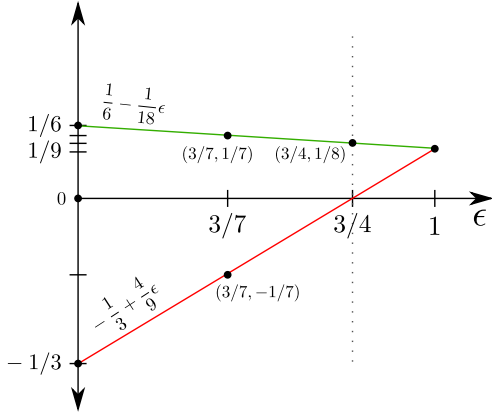


FIG. 13. **Wigner components of the noisy Strange state.** In the interval $0 \leq \epsilon < 3/7$, the negative component is larger than the positive components. At $\epsilon = 3/7$ the Wigner distribution is $(-\frac{1}{7}, \frac{1}{7}, \dots, \frac{1}{7})$. In the interval $3/7 < \epsilon < 3/4$, the positive components are larger than the negative component. In the interval $3/4 \leq \epsilon \leq 1$, there is no negativity.

Case			$m_i(n, \epsilon)$	$w_i(n, \epsilon)$
$0 \leq \epsilon < \frac{3}{7}$	n even	LHS	$8^{2i} \binom{n}{2i}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{2i} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{n-2i}$
		RHS	$8^{n-2i} \binom{n}{2i}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{n-2i} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{2i}$
	n odd	LHS	$8^{2i+1} \binom{n}{2i+1}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{2i+1} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{n-2i-1}$
		RHS	$8^{n-2i-1} \binom{n}{2i+1}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{n-2i-1} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{2i+1}$
$\frac{3}{7} \leq \epsilon < \frac{3}{4}$	n even	LHS	$8^{n-2i} \binom{n}{2i}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{n-2i} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{2i}$
		RHS	$8^{2i} \binom{n}{2i}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{2i} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{n-2i}$
	n odd	LHS	$8^{n-2i} \binom{n}{2i}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{n-2i} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{2i}$
		RHS	$8^{2i} \binom{n}{2i}$	$(\frac{1}{6} - \frac{1}{18}\epsilon)^{2i} (-\frac{1}{3} + \frac{4}{9}\epsilon)^{n-2i}$

TABLE I. Wigner components $w_i(n, \epsilon)$ of $\rho_S(\epsilon)^{\otimes n}$ along with their multiplicities $m_i(n, \epsilon)$ in decreasing order in i , $0 \leq i \leq n$. The order changes depending on the noise level ϵ , the parity of the number of copies n and the parity of the components (LHS vs RHS). Multiplication $2i$ is considered modulo $(n+1)$.

For example, the distribution of state $\rho_S(0)^{\otimes 2}$ is

$$W_{\rho_S(0)^{\otimes 2}} = \left(\overbrace{\left(-\frac{1}{3}\right)^2}^1, \overbrace{\left(\frac{1}{6}\right)^2}^{64}, \dots, \left(\frac{1}{6}\right)^2, \overbrace{-\frac{1}{3} \cdot \frac{1}{6}, \dots, -\frac{1}{3} \cdot \frac{1}{6}}^{16} \right)$$

Every standard Lorenz curve contains n elbows, labelled by

$$\{(x_i, L_i)\}_{i=-1,0,\dots,n},$$

where the boundary points $(x_{-1}, L_{-1}) = (0, 0)$ and $(x_n, L_n) = (1, 1)$ are also included. The maximum is

the $(\lfloor n/2 \rfloor)$ -th elbow and its coordinates are calculated by collecting all the positive Wigner components,

$$x_{\lfloor n/2 \rfloor} = \frac{1}{2} \left(1 + \left(\frac{7}{9} \right)^n \right), \quad (\text{E11})$$

$$L_{\lfloor n/2 \rfloor} = \frac{1}{2} \left(1 + \left(\frac{15-8\epsilon}{9} \right)^n \right). \quad (\text{E12})$$

Expressions for all the elbow coordinates follow from summing up the Wigner components in decreasing order and we present the elbow coordinates for standard Lorenz curves in Table II.

Case			x_i	L_i
$0 \leq \epsilon < \frac{3}{7}$	n even	LHS	$\Phi_+(2i; n, \frac{8}{9})$	$(\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_+(2i; n, \frac{12-4\epsilon}{15-8\epsilon})$
		RHS	$x_{\lfloor n/2 \rfloor} + \Phi_-(2i; n, \frac{1}{9})$	$L_{\lfloor n/2 \rfloor} - (\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_-(2i; n, \frac{3-4\epsilon}{15-8\epsilon})$
	n odd	LHS	$\Phi_-(2i; n, \frac{8}{9})$	$(\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_-(2i; n, \frac{12-4\epsilon}{15-8\epsilon})$
		RHS	$x_{\lfloor n/2 \rfloor} + \Phi_-(2i; n, \frac{1}{9})$	$L_{\lfloor n/2 \rfloor} - (\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_-(2i; n, \frac{3-4\epsilon}{15-8\epsilon})$
$\frac{3}{7} \leq \epsilon < \frac{3}{4}$	n even	LHS	$\Phi_+(2i; n, \frac{1}{9})$	$(\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_+(2i; n, \frac{3-4\epsilon}{15-8\epsilon})$
		RHS	$x_{\lfloor n/2 \rfloor} + \Phi_-(2i; n, \frac{8}{9})$	$L_{\lfloor n/2 \rfloor} - (\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_-(2i; n, \frac{12-4\epsilon}{15-8\epsilon})$
	n odd	LHS	$\Phi_+(2i; n, \frac{1}{9})$	$(\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_+(2i; n, \frac{3-4\epsilon}{15-8\epsilon})$
		RHS	$x_{\lfloor n/2 \rfloor} + \Phi_+(2i; n, \frac{8}{9})$	$L_{\lfloor n/2 \rfloor} - (\frac{5}{3} - \frac{8}{9}\epsilon)^n \Phi_+(2i; n, \frac{12-4\epsilon}{15-8\epsilon})$

TABLE II. Standard Lorenz curves elbow coordinates. The expression depends on the noise level ϵ , the parity of the number of copies n and the location of the elbow relative to the maximum (LHS vs RHS). Multiplication $2i$ is considered modulo $(n+1)$. Note that the Lorenz curve boundary points are $(x_{-1}, L_{-1}) := (0, 0)$ and $(x_n, L_n) = (1, 1)$.

3. Standard Lorenz curve coordinates

We can get explicit expressions for all 9^n points of the standard Lorenz curve $L_{\rho_S(\epsilon)^{\otimes n} \mathbb{1}/3}$, in terms of the elbow coordinates:

$$x_{ij} = \left(1 - \frac{j}{m_i} \right) x_{i-1} + \frac{j}{m_i} x_i, \quad (\text{E13})$$

$$L_{ij} = \left(1 - \frac{j}{m_i} \right) L_{i-1} + \frac{j}{m_i} L_i \quad (\text{E14})$$

for $j = 1, \dots, m_i$ and $i = 0, \dots, n$, where multiplicities $m_i = m_i(n, \epsilon)$ are given in Table I.

For visualisation purposes, we plot the Lorenz curve of state $\rho_S(0)^{\otimes 2}$.

Consider the state

$$\rho_S(n', \epsilon, n - n') := \rho_S(\epsilon)^{\otimes n'} \otimes \left(\frac{1}{3} \mathbb{1} \right)^{\otimes (n-n')},$$

Tensoring with the maximally mixed state keeps the Lorenz curve unchanged, but increases the resolution of

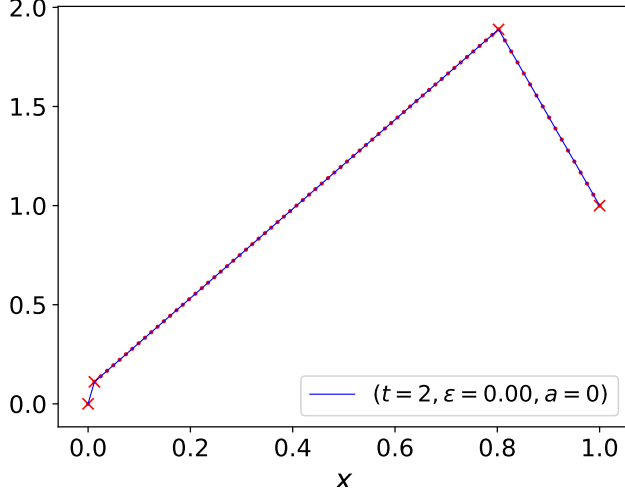


FIG. 14. **All points on the Lorenz curve are uniformly distributed.** Eqs. (E13) and (E14) capture the coordinates of all points up to the maximum.

(the uniformly distributed) points. The new point coordinates are given by:

$$x_{ijk} = (1 - p_{ijk}) x_{i-1} + p_{ijk} x_i \quad (\text{E15})$$

$$L_{ijk} = (1 - p_{ijk}) L_{i-1} + p_{ijk} L_i, \quad (\text{E16})$$

$$\text{where } p_{ijk} = \frac{k + (j-1)9^{n-n'}}{9^{n-n'} m_i}$$

for $i = 0, \dots, n'$, $j = 1, \dots, m_i(n', \epsilon')$ and $k = 1, \dots, 9^{n-n'}$.

We can unify the indices, by introducing a single index

$$I(i, j, k) := k + \left[(j-1) + \sum_{\ell=0}^{i-1} m_\ell(n', \epsilon') \right] 9^{n-n'}, \quad (\text{E17})$$

so that $I = 1, 2, \dots, 9^n$. The elbow coordinates correspond to

$$I(i, m_i(n', \epsilon'), 9^{n-n'}) = \sum_{\ell=0}^i m_\ell(n', \epsilon'), \quad i = 0, \dots, n'. \quad (\text{E18})$$

The index function I is bijective, i.e.

$$(i, j, k) = (i', j', k') \Leftrightarrow I(i, j, k) = I(i', j', k'). \quad (\text{E19})$$

4. Strange state MSD in the unital fragment

Consider the Strange state MSD process in the unital fragment,

$$\rho_S(n', \epsilon, 0) \xrightarrow{\mathcal{O}_{1/3}} \rho_S(n', \epsilon', n - n'). \quad (\text{E20})$$

We denote initial state indices without a prime and target state indices with a prime,

$$I(i, j, k = 1) = j + \sum_{\ell=0}^{i-1} m_\ell(n, \epsilon), \quad (\text{E21})$$

$$I'(i', j', k') = k' + \left[(j' - 1) + \sum_{\ell=0}^{i'-1} m_\ell(n', \epsilon') \right] 9^{n-n'}. \quad (\text{E22})$$

Pointwise Lorenz curve comparison requires $x_I = x_{I'}$, so the question is:

Given a triplet (i', j', k') , what is the tuple (i, j) such that $I(i, j) = I'(i', j', k')$?

According to Theorem 18 which is proved in Appendix D 4, for standard Lorenz curves, we need to match indices at the target state elbows, so the requirement on the indices is finding a tuple (i, j) , such that for a given $i' = 0, \dots, n'$,

$$j + \sum_{\ell=0}^{i-1} m_\ell(n, \epsilon) = \sum_{\ell=0}^{i'} m_\ell(n', \epsilon'). \quad (\text{E23})$$

As a basic example, consider the process

$$\rho_S(\epsilon)^{\otimes 4} \xrightarrow{\mathcal{O}_{1/3}} \rho_S(\epsilon')^{\otimes 2} \otimes \left(\frac{1}{3} \mathbb{1} \right)^{\otimes 2}. \quad (\text{E24})$$

We want to check which Lorenz curve is higher at the first elbows of the target state, i.e. we want to verify or reject the first inequality below:

$$L_{I(i,j)} \geq L'_{I'(0,1,81)} \quad (\text{E25})$$

where the first multiplicity of the target state is $m_0 = 1$. The multiplicities of the initial state are $(1, 384, 4096)$.

The challenge is to find i, j such that $i(i, j) = i'(0, 1, 81)$. [By trial and error], we find that $(i, j) = (1, 80)$. Now we can use Eq. (E16) to directly calculate

$$\begin{aligned} L'_{i'(0,1,81)} &= L'_0 = \left(\frac{5}{3} - \frac{8}{9}\epsilon \right)^2 \Phi_+ \left(0; 2, 4, \frac{3-\epsilon}{15-8\epsilon} \right), \\ L_{i(1,80)} &= \left(1 - \frac{80}{384} \right) L_0 + \frac{80}{384} L_1 \\ &= \left(\frac{5}{3} - \frac{8}{9}\epsilon \right)^4 \left[\frac{19}{24} \Phi_+ \left(0; 4, 4, \frac{3-\epsilon}{15-8\epsilon} \right) \right. \\ &\quad \left. + \frac{5}{24} \Phi_+ \left(2; 4, 4, \frac{3-\epsilon}{15-8\epsilon} \right) \right], \end{aligned}$$

and then compare them.

Appendix F: Technical details of bound derivation in stabiliser fragments

1. Component-multiplicity pairs

In general, a 1-copy d -dimensional state is defined exactly by its d^2 Wigner distribution W . The distribu-

tion W can always be described by a Wigner-component vector $\mathbf{w} = (w_i)_{i=1,\dots,D}$ and a multiplicity vector $\mathbf{m} = (m_i)_{i=1,\dots,D}$, where $D \leq d^2$ and they form component-multiplicity pairs denoted by $\{(w_i, m_i)\}_{i=1,\dots,D}$.

Definition 21. Consider a d^2 -dimensional distribution W and a positive integer D . We call the set of order pairs $\{(w_i, m_i)\}_{i=1,\dots,D}$ a complete set of component-multiplicity pairs, if W contains m_i components w_i and $\sum_{i=1}^D m_i = d^2$.

Therefore, such a set that describes a distribution W contains in its description each component of W exactly once. As an example, two complete sets of pairs for the Strange state are $\{(-1/3, 1), (1/6, 8)\}$ and $\{(-1/3, 1), (1/6, 2), (1/6, 3), (1/6, 3)\}$.

Consider two states ρ_A, ρ_B with Wigner distributions W_{ρ_A}, W_{ρ_B} described respectively by complete sets of component-multiplicity pairs

$$\{(w_i, m_i)\}_{i=1,\dots,D_A} \text{ and } \{(w'_j, m'_j)\}_{j=1,\dots,D_B}. \quad (\text{F1})$$

The multiplicative property of the Wigner distribution over a composite phase space $\mathcal{P}_d[d_A] \times \mathcal{P}_d[d_B]$,

$$W_{\rho_A \otimes \rho_B}(\mathbf{x}_A \oplus \mathbf{x}_B) = W_{\rho_A}(\mathbf{x}_A) W_{\rho_B}(\mathbf{x}_B), \quad (\text{F2})$$

implies that the distribution $W_{\rho_A \otimes \rho_B}$ is $d_A^2 d_B^2$ -dimensional and contains components of the form $w_i w'_j$. Therefore, the set $\{(w_i w'_j, m_i m'_j)\}$ with $i = 1, \dots, D_A$ and $j = 1, \dots, D_B$ is a complete set of component-multiplicity pairs for the distribution of the composite system $W_{\rho_A \otimes \rho_B}$. This is true because all components are of the form $w_i w'_j$ and

$$\sum_{i=1}^{D_A} \sum_{j=1}^{D_B} m_i m'_j = \sum_{i=1}^{D_A} m_i \sum_{j=1}^{D_B} m'_j = d_A^2 d_B^2.$$

Given a state ρ and a complete set of component-multiplicity pairs describing its Wigner distribution W_ρ , we now provide a method of computing the components (and multiplicities) of the n -copy distribution $W_\rho^{\otimes n}$.

Lemma 22. Let W be a d -dimensional distribution defined by a complete set of component-multiplicity pairs $\{(w_i, m_i)\}_{i=1,\dots,D}$ with $D \leq d$ and consider the distribution $W^{\otimes n}$ obtained by taking the Kronecker product $W \otimes \dots \otimes W$ between n copies of W .

Denote by $C_D^n := \{\mathbf{k}\}$ the set of all vectors $\mathbf{k} := (k_1, \dots, k_D)$ with non-negative integer components that sum to n , i.e.

$$0 \leq k_1, \dots, k_D \leq n \text{ and } k_1 + \dots + k_D = n.$$

Then, $W^{\otimes n}$ admits a complete set of component-multiplicity pairs $\{(W_{\mathbf{k}}, M_{\mathbf{k}})\}_{\mathbf{k} \in C_D^n}$, where

$$M_{\mathbf{k}} = \frac{n!}{k_1! \dots k_D!} \prod_{i=1}^D m_i^{k_i}, \quad (\text{F3})$$

$$W_{\mathbf{k}} = \prod_{i=1}^D w_i^{k_i}. \quad (\text{F4})$$

Proof. We proceed by induction.

Assume $n = 1$. Let $\mathbf{k}_i := (0, \dots, 0, 1, 0, \dots, 0)$ be a vector with its i -th component equal to 1 and 0's elsewhere. The set C_D^1 consists of all vectors of this form, i.e.

$$C_D^1 = \{\mathbf{k}_i\}_{i=1,\dots,D}$$

It is also true by direct calculation that

$$(W_{\mathbf{k}_i}, M_{\mathbf{k}_i}) = (w_i, m_i).$$

Therefore, $\{(W_{\mathbf{k}}, M_{\mathbf{k}})\}_{\mathbf{k} \in C_D^1}$ is a complete set of component-multiplicity pairs for W .

Assume that $\{(W_{\mathbf{k}}, M_{\mathbf{k}})\}_{\mathbf{k} \in C_D^n}$ as given in Eqs. (F3) and (F4) is a complete set of component-multiplicity pairs for the n -copy distribution $W^{\otimes n}$. By construction, the distribution $W^{\otimes(n+1)} = W^{\otimes n} \otimes W$ is multiplicative, so it admits the complete set of component multiplicity pairs

$$\{(W_{\mathbf{k}} w_i, M_{\mathbf{k}} m_i)\}, \mathbf{k} \in C_D^n \text{ and } i = 1, \dots, D. \quad (\text{F5})$$

Consider the component sum of the distribution $W^{\otimes(n+1)}$,

$$\begin{aligned} \sum_{\mathbf{k} \in C_D^n} \sum_{i=1}^D M_{\mathbf{k}} m_i W_{\mathbf{k}} w_i &= \sum_{\mathbf{k} \in C_D^n} M_{\mathbf{k}} W_{\mathbf{k}} \sum_{i=1}^D m_i w_i = \\ \sum_{\mathbf{k} \in C_D^n} \frac{n!}{k_1! \dots k_D!} \prod_{i=1}^D m_i^{k_i} w_i^{k_i} \sum_{i=1}^D m_i w_i &= \\ \left(\sum_{i=1}^D m_i w_i \right)^n \left(\sum_{i=1}^D m_i w_i \right) &= \left(\sum_{i=1}^D m_i w_i \right)^{n+1} = \\ \sum_{\mathbf{q} \in C_D^{n+1}} M_{\mathbf{q}} W_{\mathbf{q}}, \end{aligned}$$

where in the last expression, vectors $\mathbf{q} = (q_1, \dots, q_D)$ have non-negative integer components that sum to $(n+1)$ and

$$M_{\mathbf{q}} = \frac{(n+1)!}{q_1! \dots q_D!} \prod_{i=1}^D m_i^{q_i},$$

$$W_{\mathbf{q}} = \prod_{i=1}^D w_i^{q_i}.$$

We have used the multinomial theorem to proceed between lines 2-3 and lines 3-4 of the derivation.

We have achieved a regrouping of the distribution components. Every component $W_{\mathbf{q}}$ is of the form $W_{\mathbf{k}} w_i$ with $q_i = k_i + 1$ and $q_j = k_j$ for $j \neq i$ and

$$\sum_{\mathbf{q} \in C_D^{n+1}} M_{\mathbf{q}} = \sum_{\mathbf{q} \in C_D^{n+1}} \frac{(n+1)!}{q_1! \dots q_D!} \prod_{i=1}^D m_i^{q_i} = \left(\sum_{i=1}^D m_i \right)^{n+1} = d^{n+1},$$

which is the dimension of $W^{\otimes(n+1)}$.

Therefore, $\{(W_{\mathbf{q}}, M_{\mathbf{q}})\}_{\mathbf{q} \in C_D^{n+1}}$ is a complete set of component-multiplicity pairs for $W^{\otimes(n+1)}$, completing the proof. \square