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Approximate reversibility in the context of entropy gain, information gain, and complete positivity

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There are several inequalities in physics which limit how well we can process physical systems to achieve some intended goal, including the second law of thermodynamics, entropy bounds in quantum information theory, and the uncertainty principle of quantum mechanics. Recent results provide physically meaningful enhancements of these limiting statements, determining how well one can attempt to reverse an irreversible process. In this paper, we apply and extend these results to give strong enhancements to several entropy inequalities, having to do with entropy gain, information gain, entropic disturbance, and complete positivity of open quantum systems dynamics. Our first result is a remainder term for the entropy gain of a quantum channel. This result implies that a small increase in entropy under the action of a subunital channel is a witness to the fact that the channel's adjoint can be used as a recovery map to undo the action of the original channel. We apply this result to pure-loss, quantum-limited amplifier, and phase-insensitive quantum Gaussian channels, showing how a quantum-limited amplifier can serve as a recovery from a pure-loss channel and vice versa. Our second result regards the information gain of a quantum measurement, both without and with quantum side information. We find here that a small information gain implies that it is possible to undo the action of the original measurement if it is efficient. The result also has operational ramifications for the information-theoretic tasks known as measurement compression without and with quantum side information. Our third result shows that the loss of Holevo information caused by the action of a noisy channel on an input ensemble of quantum states is small if and only if the noise can be approximately corrected on average. We finally establish that the reduced dynamics of a system-environment interaction are approximately completely positive and trace-preserving if and only if the data processing inequality holds approximately.

Keywords: approximate reversibility, recoverability, entropy gain, information gain, entropic disturbance, divisible processes, data-processing inequality, quantum relative entropy

I. INTRODUCTION

The second law of thermodynamics constitutes a fundamental limitation on our ability to extract energy from physical systems [8, 13, 38]. The data processing inequality represents a limitation on our ability to process information, being the basis for most of the important capacity theorems in quantum information theory [47]. The entropic uncertainty principle of quantum mechanics places a limitation on how well we can measure incompatible observables [18, 24]. These seemingly disparate statements have a common mathematical foundation in an entropy inequality known as the monotonicity of quantum relative entropy [32, 43], which states that the quantum relative entropy cannot increase under the action of a quantum channel. More precisely, the quantum relative entropy between two density operators ρ and σ is defined as [44]

$$D(\rho\|\sigma) \equiv \mathrm{Tr}\{\rho\,[\log\rho - \log\sigma]\}, \tag{1}$$

and the monotonicity of quantum relative entropy states that [32, 43]

$$D(\rho\|\sigma) \geq D(\mathcal{N}(\rho)\|\mathcal{N}(\sigma)), \tag{2}$$

where $\mathcal{N}$ is a quantum channel.

Recently, researchers have explored refinements of these statements in various contexts, with the common theme being to understand how well one can attempt to reverse an irreversible process. One of the main technical developments which has allowed for these refined statements is a strengthening of the monotonicity of quantum relative entropy of the following form [48]:

$$D(\rho\|\sigma) \geq D(\mathcal{N}(\rho)\|\mathcal{N}(\sigma)) - \log F(\rho, (\mathcal{R}\circ\mathcal{N})(\rho)), \tag{3}$$

where $F(\omega,\tau) \equiv \|\sqrt{\omega}\sqrt{\tau}\|_1^2$ is the quantum fidelity [42] between two density operators ω and τ, and $\mathcal{R}$ is a recovery channel with the property that it perfectly recovers the σ state, in the sense that $\sigma = (\mathcal{R}\circ\mathcal{N})(\sigma)$ (see also [30, 40] for later developments and [19] for an important earlier development with conditional mutual information).

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Several applications follow as a consequence. Ref. [46] gave an application in thermodynamics, proving that if the free energies of two states are close and if it is possible to transition from one state to another via a thermal operation such that there is an energy gain in the process, then one can approximately reverse this thermodynamic transition without using any energy at all. Ref. [7] showed how to tighten the uncertainty principle in the presence of quantum memory [5] with another term related to how much disturbance a given measurement causes, thus unifying several aspects of quantum physics, including measurement incompatibility, entanglement, and measurement disturbance, in a single entropic uncertainty relation. Finally, Ref. [48] has given an increased understanding of many well known entropy inequalities in quantum information, such as the joint convexity of quantum relative entropy, the non-negativity of quantum discord, the Holevo bound, and multipartite information inequalities.

In this paper, we continue with this theme and derive several new results:

1. First, we give a strong improvement of the well known statement that the quantum entropy cannot decrease under the action of a unital quantum channel (a channel which preserves the identity operator). The bound that we derive has a rather simple proof, following from the operator concavity of the logarithm (related to the method used in [40]). The main physical implication of this result is that if the entropy gain under the action of a unital channel is not too large, then it is possible to reverse the action of this channel by applying its adjoint (which is a quantum channel in this case).

2. Next, we consider the information gain of a quantum measurement, a concept introduced in [22] and subsequently refined in [6, 11, 49, 50]. The information gain of a quantum measurement quantifies how much data we can gather by performing a quantum measurement on a given state. It has an operational interpretation as the rate at which a sender needs to transmit classical information to a receiver in order for them to simulate a quantum measurement on a given state [50]. Here, we prove that if the information gain is not too large, then it is possible to reverse the action of the measurement and, in the operational context, one can also simulate the measurement well on average without sending any classical data at all. The result also applies if the measurement is performed on one share of a bipartite state.

3. Third, we provide a clear operational meaning for the notion of *entropic disturbance*, defined in [12] as the loss of the Holevo information due to the action of a noisy channel on an initial ensemble of quantum states. We accomplish this by showing that a small loss of Holevo information implies that the action of the noisy channel on the input ensemble can be approximately undone, on average. This result answers a question left open from [12].

4. Finally, we give a refinement of the recent link between the data processing inequality and complete positivity of open quantum systems dynamics [9]. In [9], it was shown that the data processing inequality holds if and only if the reduced dynamics of an evolution can be described by a completely positive trace-preserving map. Here, we show how this result holds approximately, which should allow for experimental tests if desired. That is, we show that the data processing inequality holds approximately if and only if the reduced dynamics of an evolution are approximately completely positive and trace-preserving (see Section VI for precise statements).

The rest of the paper is devoted to giving more details and explanations of these results. We begin in the next section by setting notation, definitions, and reviewing the prior literature in more detail. We then follow with each of the aforementioned results and conclude in Section VII with a summary.

II. PRELIMINARIES

This section reviews background material on quantum information, all of which is available in [47]. Let $\mathcal{L}(\mathcal{H})$ denote the algebra of bounded linear operators acting on a Hilbert space $\mathcal{H}$. Let $\mathcal{L}_+(\mathcal{H})$ denote the subset of positive semi-definite operators. We also write $X \geq 0$ if $X \in \mathcal{L}_+(\mathcal{H})$. An operator ρ is in the set $\mathcal{D}(\mathcal{H})$ of density operators (or states) if $\rho \in \mathcal{L}_+(\mathcal{H})$ and $\mathrm{Tr}\{\rho\} = 1$. The tensor product of two Hilbert spaces $\mathcal{H}_A$ and $\mathcal{H}_B$ is denoted by $\mathcal{H}_A \otimes \mathcal{H}_B$ or $\mathcal{H}_{AB}$. Given a multipartite density operator $\rho_{AB} \in \mathcal{D}(\mathcal{H}_A \otimes \mathcal{H}_B)$, we unambiguously write $\rho_A = \mathrm{Tr}_B\{\rho_{AB}\}$ for the reduced density operator on system A. We use ρ_{AB}, σ_{AB}, τ_{AB}, ω_{AB}, etc. to denote general density operators in $\mathcal{D}(\mathcal{H}_A \otimes \mathcal{H}_B)$, while ψ_{AB}, φ_{AB}, ϕ_{AB}, etc. denote rank-one density operators (pure states) in $\mathcal{D}(\mathcal{H}_A \otimes \mathcal{H}_B)$ (with it implicit, clear from the context, and the above convention implying that ψ_A, φ_A, ϕ_A may be mixed if ψ_{AB}, φ_{AB}, ϕ_{AB} are pure). A purification $|\phi^\rho\rangle_{RA} \in \mathcal{H}_R \otimes \mathcal{H}_A$ of a state $\rho_A \in \mathcal{D}(\mathcal{H}_A)$ is such that $\rho_A = \mathrm{Tr}_R\{|\phi^\rho\rangle\langle\phi^\rho|_{RA}\}$. An isometry $U : \mathcal{H} \to \mathcal{H}'$ is a linear map such that $U^\dagger U = I_{\mathcal{H}}$. Often, an identity operator is implicit if we do not write it explicitly (and should be clear from the context).

Throughout this paper, we take the usual convention that $f(A) = \sum_{i:a_i\neq 0} f(a_i)|i\rangle\langle i|$ when given a function f and a Hermitian operator A with spectral decomposition $A = \sum_i a_i|i\rangle\langle i|$. In particular, A^{-1} is interpreted as a generalized inverse, so that $A^{-1} = \sum_{i:a_i\neq 0} a_i^{-1}|i\rangle\langle i|$, $\log(A) = \sum_{i:a_i>0} \log(a_i)\,|i\rangle\langle i|$, $\exp(A) = \sum_{i:a_i\neq 0} \exp(a_i)\,|i\rangle\langle i|$, etc. Throughout the paper, we interpret log as the binary logarithm. We employ the shorthand supp(A) and ker(A) to refer to the support and kernel of an operator A, respectively.

A linear map $\mathcal{N}_{A\to B} : \mathcal{L}(\mathcal{H}_A) \to \mathcal{L}(\mathcal{H}_B)$ is positive if $\mathcal{N}_{A\to B}(\sigma_A) \in \mathcal{L}(\mathcal{H}_B)_+$ whenever $\sigma_A \in \mathcal{L}(\mathcal{H}_A)_+$. Let id_A denote the identity map acting on a system A. A linear map $\mathcal{N}_{A\to B}$ is completely positive if the map $\mathrm{id}_R \otimes \mathcal{N}_{A\to B}$ is positive for a reference system R of arbitrary size. A linear map $\mathcal{N}_{A\to B}$ is trace-preserving if $\mathrm{Tr}\{\mathcal{N}_{A\to B}(\tau_A)\} = \mathrm{Tr}\{\tau_A\}$ for all input operators $\tau_A \in \mathcal{L}(\mathcal{H}_A)$. It is trace non-increasing if $\mathrm{Tr}\{\mathcal{N}_{A\to B}(\tau_A)\} \leq \mathrm{Tr}\{\tau_A\}$ for all $\tau_A \in \mathcal{L}_+(\mathcal{H}_A)$. A quantum channel is a linear map which is completely positive and trace-preserving (CPTP). A positive operator-valued measure (POVM) is a set $\{\Lambda^m\}$ of positive semi-definite operators such that $\sum_m \Lambda^m = I$. For $X, Y \in \mathcal{L}(\mathcal{H})$, let $\langle X, Y\rangle \equiv \mathrm{Tr}\{X^\dagger Y\}$ denote the Hilbert–Schmidt inner product. The adjoint $(\mathcal{M}_{A\to B})^\dagger$ of a linear map $\mathcal{M}_{A\to B}$ is the unique linear map satisfying

$$\langle Y_B, \mathcal{M}_{A\to B}(X_A)\rangle = \langle (\mathcal{M}_{A\to B})^\dagger (Y_B), X_A\rangle, \quad (4)$$

for all $X_A \in \mathcal{L}(\mathcal{H}_A)$ and $Y_B \in \mathcal{L}(\mathcal{H}_B)$. A linear map $\mathcal{M}_{A\to B}$ is unital if it preserves the identity, i.e., $\mathcal{M}_{A\to B}(I_A) = I_B$. It then follows that a linear map is unital if and only if its adjoint is trace preserving. A linear map $\mathcal{M}_{A\to B}$ is subunital if $\mathcal{M}_{A\to B}(I_A) \leq I_B$, and this is equivalent to the adjoint of $\mathcal{M}_{A\to B}$ being trace non-increasing. A quantum channel $\mathcal{U} : \mathcal{L}(\mathcal{H}_A) \to \mathcal{L}(\mathcal{H}_B)$ is an isometric channel if it has the action $\mathcal{U}(X_A) = U X_A U^\dagger$, where $X_A \in \mathcal{L}(\mathcal{H}_A)$ and $U : \mathcal{H}_A \to \mathcal{H}_B$ is an isometry.

A quantum instrument is a quantum channel that accepts a quantum system as input and outputs two systems: a classical one and a quantum one. More formally, a quantum instrument is a collection $\{\mathcal{N}^x\}$ of completely positive trace non-increasing maps, such that the sum map $\sum_x \mathcal{N}^x$ is a quantum channel. We can write the action of a quantum instrument on an input operator P as the following quantum channel:

$$P \to \sum_x \mathcal{N}^x(P) \otimes |x\rangle\langle x|, \quad (5)$$

where $\{|x\rangle\}$ is an orthonormal basis labeling the classical output of the instrument.

The trace distance between two quantum states $\rho, \sigma \in \mathcal{D}(\mathcal{H})$ is equal to $\|\rho - \sigma\|_1$. It has a direct operational interpretation in terms of the distinguishability of these states. That is, if ρ or σ are prepared with equal probability and the task is to distinguish them via some quantum measurement, then the optimal success probability in doing so is equal to $(1 + \|\rho - \sigma\|_1 /2)/2$. The fidelity is defined as $F(\rho, \sigma) \equiv \|\sqrt{\rho}\sqrt{\sigma}\|_1^2$ [42], and more generally we can use the same formula to define $F(P, Q)$ if $P, Q \in \mathcal{L}_+(\mathcal{H})$. Uhlmann's theorem states that [42]

$$F(\rho_A, \sigma_A) = \max_U |\langle \phi^\sigma|_{RA} U_R \otimes I_A |\phi^\rho\rangle_{RA}|^2, \quad (6)$$

where $|\phi^\rho\rangle_{RA}$ and $|\phi^\sigma\rangle_{RA}$ are purifications of ρ_A and σ_A, respectively, and the optimization is with respect to all isometries U_R. The same statement holds more generally for $P, Q \in \mathcal{L}_+(\mathcal{H})$. We will also use the notation $\sqrt{F}(\rho, \sigma) \equiv \|\sqrt{\rho}\sqrt{\sigma}\|_1$ to denote the "root fidelity" when convenient. The direct-sum property of the fidelity is that

$$\sqrt{F}(\omega_{XS}, \tau_{XS}) = \sum_x \sqrt{p_X(x) q_X(x)} \sqrt{F}(\omega_S^x, \tau_S^x), \quad (7)$$

for classical–quantum states

$$\omega_{XS} \equiv \sum_x p_X(x) |x\rangle\langle x|_X \otimes \omega_S^x, \quad (8)$$

$$\tau_{XS} \equiv \sum_x q_X(x) |x\rangle\langle x|_X \otimes \tau_S^x. \quad (9)$$

The quantum relative entropy $D(P\|Q)$ between $P, Q \in \mathcal{L}_+(\mathcal{H})$, with $P \neq 0$, is defined as [44]

$$D(P\|Q) = \mathrm{Tr}\{P [\log P - \log Q]\} \quad (10)$$

if $\mathrm{supp}(P) \subseteq \mathrm{supp}(Q)$ and as $+\infty$ otherwise. The relative entropy $D(P\|Q)$ is non-negative if $\mathrm{Tr}\{P\} \geq \mathrm{Tr}\{Q\}$, a result known as Klein's inequality [31]. Thus, for density operators ρ and σ, the relative entropy is non-negative, and furthermore, it is equal to zero if and only if $\rho = \sigma$. The quantum relative entropy obeys the following property:

$$D(P\|Q) \geq D(P\|Q'), \quad (11)$$

for $P, Q, Q' \in \mathcal{L}_+(\mathcal{H})$ such that $Q \leq Q'$. The following relationship between fidelity and quantum relative entropy is well known (see, e.g., [34]):

$$D(P\|Q) \geq -\log F(P, Q). \quad (12)$$

The quantum entropy $H(\rho)$ of a density operator ρ is $H(\rho) = -\mathrm{Tr}\{\rho \log \rho\}$. We often write this as $H(A)_\rho$ if ρ_A is the density operator for system A. The conditional entropy of a bipartite density operator ρ_{AB} is equal to $H(A|B)_\rho \equiv H(AB)_\rho - H(B)_\rho$. The mutual information is equal to $I(A;B)_\rho = H(A)_\rho - H(A|B)_\rho$. The conditional mutual information of a tripartite state ρ_{ABC} is equal to $I(A;B|C)_\rho = H(B|C)_\rho - H(B|AC)_\rho$. The following identities are well known (see, e.g., [47]):

$$H(A)_\rho = -D(\rho_A \| I_A), \quad (13)$$

$$H(A|B)_\rho = -D(\rho_{AB} \| I_A \otimes \rho_B), \quad (14)$$

$$I(A;B)_\rho = D(\rho_{AB} \| \rho_A \otimes \rho_B). \quad (15)$$

The following "recoverability theorem" is an enhancement of the monotonicity of quantum relative entropy (mentioned in (3)) and was proved recently in [30], by an extension of the methods from [48]:

$$D(\rho\|\sigma) \geq D(\mathcal{N}(\rho)\|\mathcal{N}(\sigma)) - \log F(\rho, (\mathcal{R}\circ\mathcal{N})(\rho)), \quad (16)$$

where $\rho \in \mathcal{D}(\mathcal{H})$, $\sigma \in \mathcal{L}_+(\mathcal{H})$, $\mathcal{N} : \mathcal{L}(\mathcal{H}) \to \mathcal{L}(\mathcal{H}')$ is a quantum channel, and $\mathcal{R}$ is a recovery quantum channel of the following form:

$$\mathcal{R}(Q) \equiv \mathrm{Tr}\{(I - \Pi_{\mathcal{N}(\sigma)})Q\}\tau + \int_{-\infty}^{\infty} dt\ p(t)\ \mathcal{R}^{t/2}_{\sigma,\mathcal{N}}(Q), \quad (17)$$

where $\Pi_{\mathcal{N}(\sigma)}$ is the projection onto the support of $\mathcal{N}(\sigma)$,

$$\tau \in \mathcal{D}(\mathcal{H}), \qquad p(t) \equiv \frac{\pi}{2}\left[\cosh(t) + 1\right]^{-1} \quad (18)$$

is a probability distribution on $t \in \mathcal{R}$,

$$\mathcal{U}_{\omega,t}(X) \equiv \omega^{it} X \omega^{-it} \quad (19)$$

for ω positive semi-definite,

$$\mathcal{P}_{\sigma,\mathcal{N}}(Q) \equiv \sigma^{1/2}\mathcal{N}^{\dagger}\left(\mathcal{N}(\sigma)^{-1/2} Q \mathcal{N}(\sigma)^{-1/2}\right)\sigma^{1/2} \quad (20)$$

is a completely positive, trace non-increasing map known as the Petz recovery map [36, 37], and $\mathcal{R}^{t}_{\sigma,\mathcal{N}}$ is a rotated or "swiveled" Petz recovery map, defined as

$$\mathcal{R}^{t}_{\sigma,\mathcal{N}} \equiv \mathcal{U}_{\sigma,-t} \circ \mathcal{P}_{\sigma,\mathcal{N}} \circ \mathcal{U}_{\mathcal{N}(\sigma),t}. \quad (21)$$

In fact, the following stronger statement holds [30]

$$D(\rho\|\sigma) \geq D(\mathcal{N}(\rho)\|\mathcal{N}(\sigma)) - \int_{-\infty}^{\infty} dt\ p(t) \log F(\rho, (\mathcal{R}^{t/2}_{\sigma,\mathcal{N}} \circ \mathcal{N})(\rho)), \quad (22)$$

which will be useful for our purposes here. The inequality in (16) implies the following one:

$$I(A;B|C)_{\rho} \geq -\log F(\rho_{ABC}, \mathcal{R}_{C\to AC}(\rho_{BC})), \quad (23)$$

where $\mathcal{R}_{C\to AC}$ is defined from (17), by taking $\sigma = \rho_{AC}$ and $\mathcal{N} = \mathrm{Tr}_A$. This follows from the definition we gave for $I(A;B|C)_{\rho}$, the equality in (14), and the inequality in (16). Similarly, the following holds as well:

$$I(A;B|C)_{\rho} \geq -\int_{-\infty}^{\infty} dt\, p(t) \log F(\rho_{ABC}, \mathcal{R}^{t/2}_{\rho_{AC},\mathrm{Tr}_A}(\rho_{BC})), \quad (24)$$

by taking $\sigma = \rho_{AC}$ and $\mathcal{N} = \mathrm{Tr}_A$. Explicitly, the action of the recovery map $\mathcal{R}^{t/2}_{\rho_{AC},\mathrm{Tr}_A}$ on an operator ω_C is given as follows:

$$\mathcal{R}^{t/2}_{\rho_{AC},\mathrm{Tr}_A}(\omega_C) = \rho_{AC}^{\frac{1-it}{2}}\left[I_A \otimes \rho_C^{-\frac{1-it}{2}}\omega_C \rho_C^{-\frac{1+it}{2}}\right]\rho_{AC}^{\frac{1+it}{2}}. \quad (25)$$

III. ENTROPY GAIN

It is well known that the quantum entropy cannot decrease under the action of a subunital, positive, and trace-preserving map [1, 2]:

$$H(\mathcal{N}(\rho)) \geq H(\rho), \quad (26)$$

where $\rho \in \mathcal{D}(\mathcal{H})$ and $\mathcal{N} : \mathcal{L}(\mathcal{H}) \to \mathcal{L}(\mathcal{H}')$ is a subunital, positive, and trace-preserving map. This entropy inequality follows as a simple consequence of the monotonicity of quantum relative entropy (now shown to hold for positive, trace-preserving maps [33]). That is, (26) follows by picking $\sigma = I$ in (2) and applying (11) and that $\mathcal{N}$ is subunital, whereby

$$\begin{aligned} -H(\rho) &= D(\rho\|I) && (27)\\ &\geq D(\mathcal{N}(\rho)\|\mathcal{N}(I)) && (28)\\ &\geq D(\mathcal{N}(\rho)\|I) && (29)\\ &= -H(\mathcal{N}(\rho)). && (30)\end{aligned}$$

This entropy inequality has a number of applications in quantum information and other contexts.

The following theorem leads to an enhancement of (26):

Theorem 1 *Let $\rho \in \mathcal{D}(\mathcal{H})$ and let $\mathcal{N} : \mathcal{L}(\mathcal{H}) \to \mathcal{L}(\mathcal{H}')$ be a positive and trace-preserving map. Then*

$$H(\mathcal{N}(\rho)) - H(\rho) \geq D(\rho\|(\mathcal{N}^{\dagger} \circ \mathcal{N})(\rho)). \quad (31)$$

Proof. This follows because

$$\begin{aligned} &H(\mathcal{N}(\rho)) - H(\rho)\\ &= \mathrm{Tr}\{\rho\log\rho\} - \mathrm{Tr}\{\mathcal{N}(\rho)\log\mathcal{N}(\rho)\} && (32)\\ &= \mathrm{Tr}\{\rho\log\rho\} - \mathrm{Tr}\{\rho\mathcal{N}^{\dagger}(\log\mathcal{N}(\rho))\} && (33)\\ &\geq \mathrm{Tr}\{\rho\log\rho\} - \mathrm{Tr}\{\rho\log(\mathcal{N}^{\dagger}\circ\mathcal{N})(\rho)\} && (34)\\ &= D(\rho\|(\mathcal{N}^{\dagger}\circ\mathcal{N})(\rho)). && (35)\end{aligned}$$

The second equality is from the definition of the adjoint. The inequality follows from operator concavity of the logarithm and the operator Jensen inequality for positive unital maps [16] (see also [40, Lemma 3.10]). ■

If $\mathcal{N}$ is additionally subunital [53], then Theorem 1 implies that $\mathcal{N}^{\dagger}$ is trace non-increasing, which in turn implies that $D(\rho\|(\mathcal{N}^{\dagger}\circ\mathcal{N})(\rho)) \geq 0$ by Klein's inequality. Thus, in this case, we obtain a significant strengthening of the well known fact that the entropy increases under the action of a subunital, positive, trace-preserving map.

The resulting entropy inequality also leads to an interpretation in terms of recoverability, in the sense discussed in [48]. That is, we can take the recovery map to be

$$\mathcal{R}(Y) \equiv \mathcal{N}^{\dagger}(Y) + \mathrm{Tr}\{(\mathrm{id} - \mathcal{N}^{\dagger})(Y)\}\tau, \quad (36)$$

where τ is any state in $\mathcal{D}(\mathcal{H})$, and we get that, for all ρ,

$$H(\mathcal{N}(\rho)) - H(\rho) \geq D(\rho\|(\mathcal{R}\circ\mathcal{N})(\rho)) \quad (37)$$

by applying (11), because $(\mathcal{R}\circ\mathcal{N})(\rho) \geq (\mathcal{N}^{\dagger}\circ\mathcal{N})(\rho)$. Note that $\mathcal{R}$ is a positive map if $\mathcal{N}$ is. We also note that if $\mathcal{N}$ is subunital the quantity $D(\rho\|(\mathcal{R}\circ\mathcal{N})(\rho))$ can be viewed as a measure of how much $\mathcal{N}$ deviates from being an isometry, being equal to zero if $\mathcal{N}$ is an isometric channel and non-zero otherwise (here we could also maximize the quantity with respect to input states ρ and output states τ).

Thus, what we find is an improvement over what we would get by applying (16) or the main result of [40]. First, there is a mathematical advantage in the sense that $\mathcal{N}$ is not required to be a channel, but it suffices for it to be a positive map. This addresses an open question from [33] for a very special case. Some might also consider this to be a physical advantage as well, given that in some situations the description of quantum dynamical evolutions is not given by a completely positive map (see, e.g., [9] and references therein). Second, the remainder term in (37) features the quantum relative entropy and thus is stronger than the $-\log F$ bound in (16) (cf. (12)) and the "measured relative entropy" term from [40]. Finally, note that Theorem 1 represents an improvement of some of the results from [15, 52].

A. Application to bosonic channels

Theorem 1 finds application for practical bosonic channels that have a long history in quantum information theory, in particular, the pure-loss and quantum-limited amplifier channels, and even all phase insensitive Gaussian channels [45]. A pure-loss channel is defined from the following input-output Heisenberg-picture relation:

$$\hat{b} = \sqrt{\eta}\hat{a} + \sqrt{1-\eta}\hat{e}, \tag{38}$$

where $\hat{a}$, $\hat{b}$, and $\hat{e}$ are the field-mode annihilation operators representing the sender's input, the receiver's output, and the environmental input of the channel. The parameter $\eta \in [0,1]$ represents the average fraction of photons that make it from the sender to receiver. For the pure-loss channel, the environment is prepared in the vacuum state. Let $\mathcal{B}_\eta$ denote the CPTP map corresponding to this channel. A quantum-limited amplifier channel is defined from the following input-output Heisenberg-picture relation:

$$\hat{b} = \sqrt{G}\hat{a} + \sqrt{G-1}\hat{e}^\dagger, \tag{39}$$

where $\hat{a}$, $\hat{b}$, and $\hat{e}$ have the same physical meaning as given for the pure-loss channel. The parameter $G \in [1,\infty)$ represents the gain or amplification factor of the channel. For the quantum-limited amplifier channel, the environment is prepared in the vacuum state. Let $\mathcal{A}_G$ denote the CPTP map corresponding to this channel.

One of the critical insights of [28] is that these channels are "almost unital," in the sense that

$$\mathcal{B}_\eta(I) = \eta^{-1}I, \qquad \mathcal{A}_G(I) = G^{-1}I, \tag{40}$$

and that their adjoints are strongly related, in the sense that

$$\mathcal{B}_\eta^\dagger = \eta^{-1}\mathcal{A}_{1/\eta}, \tag{41}$$

$$\mathcal{A}_G^\dagger = G^{-1}\mathcal{B}_{1/G}. \tag{42}$$

Observe that the pure-loss channel is superunital and the amplifier channel is subunital. These facts allow us to apply Theorem 1 and the fact that $D(\rho\|c\sigma) = D(\rho\|\sigma) - \log c$ for $c > 0$ to find that

$$H(\mathcal{B}_\eta(\rho)) - H(\rho) \geq D(\rho\|(\mathcal{A}_{\frac{1}{\eta}} \circ \mathcal{B}_\eta)(\rho)) + \log\eta, \tag{43}$$

$$H(\mathcal{A}_G(\rho)) - H(\rho) \geq D(\rho\|(\mathcal{B}_{\frac{1}{G}} \circ \mathcal{A}_G)(\rho)) + \log G. \tag{44}$$

These bounds demonstrate that a quantum-limited amplifier suffices as a reversal channel for a pure-loss channel and vice versa. Note that the above reversal is only good for weak losses and weak amplifiers (i.e., if $\eta \approx 1$ or $G \approx 1$). We can also conclude that

$$H((\mathcal{A}_G \circ \mathcal{B}_\eta)(\rho)) - H(\rho) \geq \\ D(\rho\|(\mathcal{A}_{1/\eta} \circ \mathcal{B}_{1/G} \circ \mathcal{A}_G \circ \mathcal{B}_\eta)(\rho)) + \log[\eta G], \tag{45}$$

because

$$(\mathcal{A}_G \circ \mathcal{B}_\eta)^\dagger = [\eta G]^{-1} \mathcal{A}_{1/\eta} \circ \mathcal{B}_{1/G}. \tag{46}$$

The above bound applies to any phase insensitive quantum Gaussian channel, given that any such channel can be written as a serial concatenation of a pure-loss channel and a quantum-limited amplifier channel [14, 21].

B. Optimized entropy gain

In [3], the minimal entropy gain of a quantum channel was defined as

$$G(\mathcal{N}) \equiv \inf_\rho [H(\mathcal{N}(\rho)) - H(\rho)], \tag{47}$$

and the following bounds were established for a channel with the same input and output Hilbert space $\mathcal{H}$:

$$-\log\dim(\mathcal{H}) \leq G(\mathcal{N}) \leq 0. \tag{48}$$

(See also [25–27] for related work.) Applying Theorem 1 gives the following alternate lower bound for the entropy gain of a quantum channel:

$$G(\mathcal{N}) \geq \inf_\rho D(\rho\|(\mathcal{N}^\dagger \circ \mathcal{N})(\rho)). \tag{49}$$

C. Entropy gain in the presence of quantum side information

A generalization of the entropy inequality in (37) holds for the case of the conditional quantum entropy, found by applying the same method:

Corollary 2 *Let $\rho_{AB} \in \mathcal{D}(\mathcal{H}_A \otimes \mathcal{H}_B)$ and $\mathcal{N}_{A\to A'} \otimes \mathrm{id}_B : \mathcal{L}(\mathcal{H}_{AB}) \to \mathcal{L}(\mathcal{H}_{A'B})$ be a positive and trace-preserving map. Then*

$$H(A'|B)_\sigma - H(A|B)_\rho \\ \geq D(\rho_{AB}\|((\mathcal{N}_{A\to A'})^\dagger \circ \mathcal{N}_{A\to A'})(\rho_{AB})), \tag{50}$$

where $\sigma_{A'B} \equiv (\mathcal{N}_{A\to A'} \otimes \mathrm{id}_B)(\rho_{AB})$.

Proof. This follows by applying Theorem 1 and definitions. From Theorem 1, we can conclude that

$$H(A'B)_\sigma - H(AB)_\rho \geq D(\rho_{AB}\|((\mathcal{N}_{A\to A'})^\dagger \circ \mathcal{N}_{A\to A'})(\rho_{AB})). \quad (51)$$

Consider also that

$$H(A'B)_\sigma - H(AB)_\rho$$
$$= H(A'B)_\sigma - H(B)_\sigma - [H(AB)_\rho - H(B)_\rho] \quad (52)$$
$$= H(A'|B)_\sigma - H(A|B)_\rho, \quad (53)$$

where we have used that $H(B)_\rho = H(B)_\sigma$. Combining these gives (50). ■

Remark 3 *In the above corollary, note that we need not necessarily take the map $\mathcal{N}_{A\to A'}$ to be completely positive—we merely require that $\mathcal{N}_{A\to A'} \otimes \mathrm{id}_B$ be a positive map. For example, if system B is a qubit, then we only require $\mathcal{N}_{A\to A'}$ to be two-positive in order for the corollary to apply.*

IV. INFORMATION GAIN

Groenewold originally defined the information gain of a quantum instrument $\{\mathcal{N}^x\}$, when performed on a quantum state ρ_A, as follows [22]:

$$I_G(\{\mathcal{N}^x\}, \rho_A) \equiv H(\rho_A) - \sum_x p_X(x) H(\rho^x_{A'}), \quad (54)$$

where

$$\rho^x_{A'} \equiv \frac{\mathcal{N}^x_{A\to A'}(\rho_A)}{p_X(x)}, \qquad p_X(x) \equiv \mathrm{Tr}\{\mathcal{N}^x_{A\to A'}(\rho_A)\}. \quad (55)$$

This definition was based on the physical intuition that information gain should be identified with the entropy reduction of the measurement. However, it was later realized that the entropy reduction can be negative, and that this happens if and only if the instrument is not an efficient measurement (an efficient measurement is such that each $\mathcal{N}^x$ consists of a single Kraus operator [20, 35]).

Apparently without realizing the connection to Groenewold's information gain of a measurement, Winter considered the operational, information-theoretic task [50] of determining the rate at which classical information would need to be communicated from a sender to a receiver in order to simulate the action of the measurement on a given state (if shared randomness is allowed for free between sender and receiver). He called this task "measurement compression," given that the goal is to send the classical output of the measurement at the smallest rate possible, in such a way that a third party would not be able to distinguish the output of the protocol performed on many copies of ϕ^ρ_{RA} from the same number of copies of the following state:

$$\sigma_{RX} \equiv \sum_x \mathrm{Tr}_{A'}\{(\mathrm{id}_R \otimes \mathcal{N}^x_{A\to A'})(\phi^\rho_{RA})\} \otimes |x\rangle\langle x|_X, \quad (56)$$

where ϕ^ρ_{RA} is a purification of ρ and $\{|x\rangle\}$ is an orthonormal basis for the classical output X of the measurement. He found that the optimal rate of measurement compression is equal to the mutual information of the measurement $I(R;X)_\sigma$.

After Winter's development, Ref. [11] suggested that the information gain of the measurement should be defined as its mutual information. The advantage of such an approach is that the mutual information $I(R;X)_\sigma$ is non-negative and has a clear operational interpretation. Furthermore, it is equal to the entropy reduction in (54) for efficient measurements [11] and thus connects with Groenewold's original intuition.

Winter's result was later extended in two different directions. First, Ref. [49] allowed for a correlated initial state ρ_{AB}, shared between the sender and receiver before communication begins. In this case, the optimal rate at which the sender needs to transmit classical information in order to simulate the measurement is equal to the conditional mutual information $I(R;X|B)_\omega$, where the conditional mutual information is with respect to the following state:

$$\omega_{RBX} \equiv \sum_x \mathrm{Tr}_{A'}\{(\mathrm{id}_R \otimes \mathcal{N}^x_{A\to A'})(\phi^\rho_{RAB})\} \otimes |x\rangle\langle x|_X, \quad (57)$$

and ϕ^ρ_{RAB} is a purification of ρ_{AB}. We can thus call $I(R;X|B)_\omega$ the information gain in the presence of quantum side information (IG-QSI), and the information-processing task is known as measurement compression with quantum side information [49]. In general, the IG-QSI is smaller than $I(RB;X)_\omega$, which is the rate at which classical communication would need to be transmitted if the receiver does not make use of the B system. The other extension of Winter's result was to determine the rate required to simulate the instrument on an arbitrary input state, and the optimal rate was proved to be equal to the optimized information gain

$$\max_\rho I(R;X)_\sigma, \quad (58)$$

where the optimization is with respect to all input states ρ_A leading to a purification ϕ^ρ_{RA} [6].

A. General bounds on the information gain

Let us consider now the channel $\mathcal{N}$ associated to a quantum instrument $\{\mathcal{N}^x\}$, as defined in (5). By defining the state $\sigma_{A'X}$ as

$$\sigma_{A'X} \equiv \mathcal{N}_{A\to A'X}(\rho_A) \quad (59)$$
$$= \sum_x \mathcal{N}^x_{A\to A'}(\rho_A) \otimes |x\rangle\langle x|_X \quad (60)$$
$$\equiv \sum_x p_X(x) \rho^x_{A'} \otimes |x\rangle\langle x|_X, \quad (61)$$

Theorem 1 in this case gives

$$H(A'X)_\sigma - H(A)_\rho$$
$$= H(X)_\sigma + \sum_x p_X(x) H(\rho^x_{A'}) - H(\rho_A) \quad (62)$$
$$= H(X)_\sigma - I_G(\{\mathcal{N}^x\}, \rho_A) \quad (63)$$
$$\geq D(\rho_A \| (\mathcal{N}^\dagger \circ \mathcal{N})(\rho_A)), \quad (64)$$

namely,

$$I_G(\{\mathcal{N}^x\}, \rho_A) \leq H(X)_\sigma - D(\rho_A \| (\mathcal{N}^\dagger \circ \mathcal{N})(\rho_A)). \quad (65)$$

The above upper bound on Groenewold's information gain is valid for any quantum instrument $\{\mathcal{N}^x\}$ and any state ρ.

A much tighter bound can be given if the instrument is efficient. In this case, it is easy to prove that the channel $\mathcal{N}$ defined in (5) is always subunital. This a consequence of the fact that, if $\mathcal{N}^x_{A\to A'}(C_A) = V^x_{A\to A'} \sqrt{\Lambda^x_A} C_A \sqrt{\Lambda^x_A} (V^x_{A\to A'})^\dagger$, where Λ^x_A are POVM elements and $V^x_{A\to A'}$ are isometries, then $\mathcal{N}^X_{A\to A'}(\mathbb{1}_A) = V^x_{A\to A'} \Lambda^x_A (V^x_{A\to A'})^\dagger \leq \mathbb{1}_{A'}$, for all x. Moreover, for efficient measurements Groenewold's information gain and the mutual information of the measurement $I(R;X)_\sigma$ are equal [11].

Thus, for efficient quantum instruments, Theorem 1 leads to the following bound:

$$H(X)_\sigma - I(R;X)_\sigma = H(X|R)_\sigma \quad (66)$$
$$\geq D(\rho_A \| (\mathcal{R} \circ \mathcal{N})(\rho_A)). \quad (67)$$

where $\mathcal{R}$ is a recovery channel independent of ρ. In other words, whenever the reference R and the classical outcome X are almost perfectly correlated, i.e., $H(X|R)_\sigma \approx 0$, then the action of the instrument on the input state ρ can be almost perfectly corrected on average, i.e., $D(\rho_A \| (\mathcal{R} \circ \mathcal{N})(\rho_A)) \approx 0$.

We notice here that the quantity in (66) has been given an interesting thermodynamical interpretation in Ref. [29], so that the above bound can be seen as a strengthening of the second law for efficient quantum measurements.

The above bound also provides a way to quantify, in an information-theoretic way, "how close" a given POVM is to the ideal measurement of an observable: one just need to prepare a state ρ that commutes with that observable (for example, the maximally mixed state $\mathbb{1}/d$) and feed it through an efficient measurement of the given POVM. The entropy difference in (66) is then a good indicator of such a "closeness," being null whenever the POVM corresponds to a sharp measurement along the diagonalizing basis. This method is somewhat similar to the approach introduced in Ref. [10].

B. Information gain without quantum side information

In what follows, we demonstrate how the refined entropy inequalities in (22) and (24) have implications for the information gain of a quantum measurement, both without and with quantum side information. We begin with the simpler case of information gain without quantum side information, a scenario considered in [11]. The theorem below gives a lower bound on the information gain in terms of how well one can recover from the action of an efficient measurement. It can be viewed as a corollary of the more general statement given in Theorem 5 in the next section.

Theorem 4 *Let $\rho \in \mathcal{D}(\mathcal{H}_A)$ and $\{\mathcal{N}^x\}$ be a quantum instrument, where each $\mathcal{N}^x : \mathcal{L}(\mathcal{H}_A) \to \mathcal{L}(\mathcal{H}_{A'})$. Then the following inequality holds*

$$I(R;X)_\sigma \geq -\log F(\sigma_{RX}, \sigma_R \otimes \sigma_X). \quad (68)$$

If the quantum instrument is efficient, then the above inequality implies that

$$I(R;X)_\sigma \geq -2\log\left[\sum_x p_X(x) \sqrt{F}(\mathcal{U}^x_{A'\to A}(\phi^{\rho_x}_{RA'}), \phi^\rho_{RA})\right], \quad (69)$$

for some collection $\{\mathcal{U}^x_{A'\to A}\}$, where each $\mathcal{U}^x_{A'\to A}$ is an isometric quantum channel, $\phi^{\rho_x}_{RA'}$ is a purification of $\rho^x_{A'}$ defined in (55)*, and $p_X(x)$ is defined in* (55)*.*

Proof. The inequality in (68) is a simple consequence of (15) and (12). The inequality in (69) follows because

$$\sqrt{F}(\sigma_{RX}, \sigma_R \otimes \sigma_X) = \sum_x p_X(x) \sqrt{F}(\phi^{\rho_x}_R, \phi^\rho_R). \quad (70)$$

Applying Uhlmann's theorem (see (6)), we can conclude that there exist isometric channels $\mathcal{U}^x_{A'\to A}$ such that $F(\phi^{\rho_x}_R, \phi^\rho_R) = F(\mathcal{U}^x_{A'\to A}(\phi^{\rho_x}_{RA'}), \phi^\rho_{RA})$ for all x. ■

The implication of the inequality in (69) is that if the information gain of the measurement is small, so that

$$I(R;X)_\sigma \approx 0, \quad (71)$$

then it is possible to reverse the action of the measurement approximately, in such a way as to restore the post-measurement state to the original state with a fidelity

$$\sum_x p_X(x) \sqrt{F}(\mathcal{U}^x_{A'\to A}(\phi^{\rho_x}_{RA'}), \phi^\rho_{RA}) \approx 1. \quad (72)$$

We can thus view this result as a one-sided information-disturbance trade-off. Note that [11, Theorem 1] contains an observation related to this. The observation above is also related to the general one from [39], but the result above is stronger: an inability to find correction isometries, which leads to a small fidelity, is a witness to having a large information gain $I(R;X)_\sigma$, due to the presence of the negative logarithm in (69).

The inequality in (69) also has an operational implication for Winter's measurement compression task. If the information gain is small, so that (71) holds, then the sender and receiver can simulate the measurement with a high fidelity per copy of the source state, in such a way that the sender does not need to transmit any classical information at all. The receiver can just prepare many copies of ρ_A locally, perform the measurements, and deliver the outputs of the measurements as the classical data. This situation occurs because the reference system R is approximately decoupled from the classical output, in the sense that $F(\sigma_{RX}, \sigma_R \otimes \sigma_X) \approx 1$ if $I(R;X)_\sigma \approx 0$.

C. Information gain with quantum side information

We can readily extend the above results to the case of quantum side information, by employing the inequality in (23). This leads to the following theorem:

Theorem 5 *Let $\rho_{AB} \in \mathcal{D}(\mathcal{H}_A \otimes \mathcal{H}_B)$ and $\{\mathcal{N}^x\}$ be a quantum instrument, where each $\mathcal{N}^x : \mathcal{L}(\mathcal{H}_A) \to \mathcal{L}(\mathcal{H}_{A'})$. Then the following inequality holds*

$$I(R;X|B)_\omega \geq -2 \int_{-\infty}^{\infty} dt\ p(t) \log \left[\sum_x p_X(x) \sqrt{F}(\omega^x_{RB}, \mathcal{R}^{x,t/2}_B(\omega_{RB})) \right], \tag{73}$$

where

$$\omega_{RBX} \equiv \sum_x \mathrm{Tr}_{A'}\{\mathcal{N}^x_{A\to A'}(\phi^\rho_{RAB})\} \otimes |x\rangle\langle x|_X, \tag{74}$$

ϕ^ρ_{RAB} is a purification of ρ_{AB},

$$\omega^x_{RBA'} \equiv \frac{\mathcal{N}^x_{A\to A'}(\phi^\rho_{RAB})}{p_X(x)}, \tag{75}$$

$$p_X(x) \equiv \mathrm{Tr}\{\mathcal{N}^x_{A\to A'}(\phi^\rho_{RAB})\}, \tag{76}$$

$\{p_X(x)\mathcal{R}^{x,t/2}_B\}$ is a quantum instrument defined by

$$\mathcal{R}^{x,t/2}_B(\omega_{RB}) \equiv (\omega^x_B)^{\frac{1-it}{2}}\, \omega_B^{-\frac{1-it}{2}} (\omega_{RB}) \omega_B^{-\frac{1+it}{2}} (\omega^x_B)^{\frac{1+it}{2}}, \tag{77}$$

and $p(t)$ is defined in (18). If the instrument $\{\mathcal{N}^x\}$ is efficient, then the following inequality holds as well:

$$I(R;X|B)_\omega \geq -2 \int_{-\infty}^{\infty} dt\ p(t) \log \left[\sum_x p_X(x) \sqrt{F_{x,t}} \right], \tag{78}$$

for some collection $\{\mathcal{U}^{x,t}_{A\to A'}\}$, where

$$F_{x,t} \equiv F(\omega^x_{RBA'}, (\mathcal{R}^{x,t/2}_B \otimes \mathcal{U}^{x,t}_{A\to A'})(\phi^\rho_{RBA})) \tag{79}$$

and each $\mathcal{U}^{x,t}_{A\to A'}$ is an isometric quantum channel.

Proof. We begin by proving the inequality in (73). Consider that

$$I(R;X|B)_\omega \geq - \int_{-\infty}^{\infty} dt\ p(t) \log F(\omega_{RBX}, \mathcal{R}^{t/2}_{B\to BX}(\omega_{RB})), \tag{80}$$

where

$$\mathcal{R}^{t/2}_{B\to BX}(\omega_{RB}) = \omega_{BX}^{\frac{1-it}{2}} \omega_B^{-\frac{1-it}{2}} \omega_{RB} \omega_B^{-\frac{1+it}{2}} \omega_{BX}^{\frac{1+it}{2}}, \tag{81}$$

which is a direct consequence of (24). By a direct calculation, we find that

$$\mathcal{R}^{t/2}_{B\to BX}(\omega_{RB}) = \sum_x p_X(x)|x\rangle\langle x|_X \otimes \mathcal{R}^{x,t/2}_B(\omega_{RB}), \tag{82}$$

with $\mathcal{R}^{x,t/2}_B$ defined in (77). This then leads to the inequality in (73), by applying the direct sum property of fidelity. The inequality in (78) is an application of Uhlmann's theorem, after observing that the rank-one operator $\mathcal{R}^{x,t/2}_B(\phi^\rho_{RBA})$ purifies $\mathcal{R}^{x,t/2}_B(\omega_{RB})$ and the rank-one operator $\omega^x_{RBA'}$ purifies ω^x_{RB}. The aforementioned operators are rank-one if the measurement is efficient (which is what we assumed in the statement of the theorem). ■

The implications of Theorem 5 are similar to those of Theorem 4, except they apply to a setting in which quantum side information is available. If the information gain of the measurement is small, so that

$$I(R;X|B)_\omega \approx 0, \tag{83}$$

then it is possible to reverse the action of the measurement approximately, in such a way as to restore the post-measurement state of systems RA' to the original state on systems RA with a fidelity larger than

$$\int_{-\infty}^{\infty} dt\ p(t) \sum_x p_X(x) \sqrt{F_{x,t}} \approx 1. \tag{84}$$

This follows from the concavity of the fidelity. The reversal operation consists of two steps. First, Bob performs the instrument $\{p_X(x)\mathcal{R}^{x,t/2}_B\}$. He then forwards the outcomes to Alice, who performs a channel corresponding to the inverse of the isometric quantum channel $\mathcal{U}^{x,t}_{A\to A'}$. Then, the average fidelity is high if the information gain is small. We can view this result as a one-sided information-disturbance trade-off which extends the aforementioned one without quantum side information.

The inequality in (78) also has an operational implication for measurement compression with quantum side information [49]. If the IG-QSI is small, so that (83) holds, then the sender and receiver can simulate the measurement with a high fidelity per copy of the source state, in such a way that the sender does not need to transmit any classical information at all. The receiver can just perform the instrument $\{p_X(x)\mathcal{R}^{x,t/2}_B\}$ with probability $p(t)$ on the individual B systems of many copies of ρ_{AB} and deliver the classical outputs of the measurements as the classical data. This situation occurs because the X system of ω_{RBX} is approximately recoverable from B alone, in the sense that $\int_{-\infty}^{\infty} dt\ p(t) \sum_x p_X(x)\sqrt{F_{x,t}} \approx 1$ if $I(R;X|B)_\omega \approx 0$. This latter result might have implications for quantum communication complexity (cf. [41]).

V. ENTROPIC DISTURBANCE

Ref. [12] (see in particular Section 5 therein) considered the possibility of introducing an entropic measure of average disturbance as follows. Imagine that an initial ensemble of quantum states $\mathcal{E} = \{p_X(x); \rho_A^x\}_x$ is fed through a quantum channel $\mathcal{N} : \mathcal{L}(\mathcal{H}_A) \to \mathcal{L}(\mathcal{H}_{A'})$. Consider the Holevo information of the initial ensemble $\mathcal{E}$:

$$\chi(\mathcal{E}) \equiv H(\rho_A^{\mathcal{E}}) - \sum_x p_X(x) H(\rho_A^x), \tag{85}$$

where $\rho_A^{\mathcal{E}}$ denotes the average quantum state $\sum_x p_X(x)\rho_A^x$. By the monotonicity of the Holevo information, the following inequality holds

$$\Delta\chi(\mathcal{E}) \equiv \chi(\mathcal{E}) - \chi(\mathcal{N}(\mathcal{E})) \geq 0, \tag{86}$$

where by $\mathcal{N}(\mathcal{E})$ we mean the output ensemble $\{p_X(x); \mathcal{N}(\rho_{A'}^x)\}_x$.

It is known that the condition $\Delta\chi(\mathcal{E}) = 0$ implies the existence of a recovery CPTP linear map $\mathcal{R} : \mathcal{L}(\mathcal{H}_{A'}) \to \mathcal{L}(\mathcal{H}_A)$ such that

$$\mathcal{R} \circ \mathcal{N}(\rho_A^x) = \rho_A^x, \tag{87}$$

for all x [23]. In Ref. [12] the question was considered, whether a similar conclusion would hold also in the approximate case, but an answer was given only in the case in which the input ensemble consists of two mutually unbiased bases distributed with uniform prior, as done in [17].

Recent results about approximate recoverability give a solution to this problem, by demonstrating that there exists a recovery channel that can approximately recover if the loss of Holevo information is small. In fact, a special case of this problem was already solved in [48, Corollary 16] when $\mathcal{N}$ is a measurement channel. Here we establish the following more general theorem:

Theorem 6 *Let $\mathcal{E} = \{p_X(x), \rho_A^x\}$ be an ensemble of states in $\mathcal{D}(\mathcal{H}_A)$ and $\mathcal{N} : \mathcal{L}(\mathcal{H}_A) \to \mathcal{L}(\mathcal{H}_{A'})$ a quantum channel. Then there exists a recovery channel $\mathcal{R} : \mathcal{L}(\mathcal{H}_{A'}) \to \mathcal{L}(\mathcal{H}_A)$ such that*

$$\chi(\mathcal{E}) - \chi(\mathcal{N}(\mathcal{E})) \geq -2\log \sum_x p_X(x)\sqrt{F}(\rho_A^x, (\mathcal{R}\circ\mathcal{N})(\rho_A^x)). \tag{88}$$

Proof. Introduce an auxiliary system X and the bipartite classical-quantum state

$$\rho_{XA} \equiv \sum_x p_X(x)|x\rangle\langle x|_X \otimes \rho_A^x, \tag{89}$$

where the vectors $\{|x\rangle\}$ are orthonormal in the Hilbert space $\mathcal{H}_X$. In an analogous way, we also write

$$\sigma_{XA'} \equiv \sum_x p_X(x)|x\rangle\langle x|_X \otimes \mathcal{N}_{A\to A'}(\rho_A^x). \tag{90}$$

Then,

$$\begin{aligned}
&\chi(\mathcal{E}) - \chi(\mathcal{N}(\mathcal{E})) \\
&= I(X;A)_\rho - I(X;A')_\sigma && (91)\\
&= -H(X|A)_\rho + H(X|A')_\sigma && (92)\\
&= D(\rho_{XA}\|I_X\otimes\rho_A) \\
&\quad - D((\mathrm{id}_X\otimes\mathcal{N}_{A\to A'})(\rho_{XA})\|(\mathrm{id}_X\otimes\mathcal{N}_{A\to A'})(I_X\otimes\rho_A)). && (93)
\end{aligned}$$

We now invoke (16), noticing that, in this case, the operator σ has the special form $I_X \otimes \rho_A$ and the noise acts only locally, i.e., it has the form $\mathrm{id}_X \otimes\mathcal{N}_{A\to A'}$. These two facts together imply that (16) can be written in this case as follows:

$$\chi(\mathcal{E}) - \chi(\mathcal{N}(\mathcal{E})) \geq -\log F(\rho_{XA}, (\mathrm{id}_X\otimes\mathcal{R}_{A'\to A})(\sigma_{XA'})), \tag{94}$$

where $\mathcal{R} : \mathcal{L}(\mathcal{H}_{A'}) \to \mathcal{L}(\mathcal{H}_A)$ is a suitable recovery channel. Finally, we make use of the direct sum property of the fidelity in (7) to establish (88). ■

VI. COMPLETELY POSITIVE TRACE-PRESERVING MAPS AND QUANTUM DATA PROCESSING

This section demonstrates how the inequality in (23) and the Alicki–Fannes–Winter inequality [4, 51] lead to a robust version of the main conclusion of [9], which links the data processing inequality to complete positivity of open quantum systems dynamics.

There, the problem of open quantum systems evolution in the presence of initial system-environment correlations was considered. In fact, if the system and its surrounding environment are correlated already before the interaction governing their joint evolution is turned on, then in general there does not necessarily exist a linear (let alone positive or even completely positive) map describing the reduced dynamics of the system [8]. Ref. [9] proposed to use the quantum data-processing inequality as a criterion to establish whether the system's reduced dynamics are compatible with a CPTP linear map or not.

The operational framework considered in [9] can be summarized as follows:

1. It is assumed that possible joint system-environment states belong to a known family of states that constitutes the promise to the problem. It is also assumed that such a family is "steerable," namely, that there exists a tripartite density operator ρ_{RQE} such that the reference R is able to steer all possible bipartite system-environment states in the family. Such a condition encompasses essentially all cases considered in the literature. We therefore assume that, at some initial time $t = \tau$, the system-environment correlations can be described by means of *one* given tripartite state ρ_{RQE}.

2. Moving to the next instant in time, $t = \tau + \Delta$, the system-environment pair has evolved according to some isometry $V : QE \to Q'E'$, while the reference R remains unchanged. The tripartite configuration ρ_{RQE} has correspondingly evolved to the tripartite configuration $\sigma_{RQ'E'} = (I_R \otimes V_{QE})\rho_{RQE}(I_R \otimes V_{QE}^\dagger)$.

3. Only at this point we focus on the reduced reference-system dynamics (i.e., the transformation mapping ρ_{RQ} to $\sigma_{RQ'}$), checking whether these are compatible with the application of a CPTP linear map on the system Q *alone*. More explicitly, we check whether there exists a CPTP linear map $\mathcal{E} : Q \to Q'$ such that $\sigma_{RQ'} = (\mathrm{id}_R \otimes \mathcal{E}_Q)(\rho_{RQ})$.

The following theorem generalizes to the approximate scenario the insight provided in Ref. [9].

Theorem 7 *Fix a tripartite configuration ρ_{RQE}. Suppose that the data processing inequality holds approximately for all joint system-environment evolutions $V_{QE\to Q'E'}$, i.e.,*

$$I(R;Q')_\sigma \leq I(R;Q)_\rho + \varepsilon, \tag{95}$$

where $\varepsilon > 0$ and

$$\sigma_{RQ'E'} = V_{QE\to Q'E'}\rho_{RQE}V^\dagger_{QE\to Q'E'}. \tag{96}$$

Then the conditional mutual information is nearly equal to zero:

$$I(R;E|Q)_\rho \leq \varepsilon, \tag{97}$$

and the reduced dynamics are approximately CPTP, i.e., to every unitary interaction $V_{QE\to Q'E'}$, there exists a CPTP map $\mathcal{E}_{Q\to Q'}$ such that

$$-\log F(\sigma_{RQ'}, \mathcal{E}_{Q\to Q'}(\rho_{RQ})) \leq \varepsilon. \tag{98}$$

Proof. We begin by proving (97) with the same approach used in [9]. Consider the particular evolution in which $Q' = QE$ and system E' is trivial. The assumption that data processing holds approximately gives that

$$I(R;Q)_\rho + \varepsilon \geq I(R;Q')_\sigma = I(R;QE)_\rho. \tag{99}$$

We can rewrite this inequality using the chain rule for conditional mutual information as

$$\varepsilon \geq I(R;QE)_\rho - I(R;Q)_\rho = I(R;E|Q)_\rho, \tag{100}$$

which proves (97). Now, from the inequality in (23), we know that there exists a recovery map $\mathcal{R}_{Q\to QE}$ such that

$$I(R;E|Q)_\rho \geq -\log F\left(\rho_{RQE}, \mathcal{R}_{Q\to QE}\left(\rho_{RQ}\right)\right) \tag{101}$$

Since the fidelity is invariant with respect to unitaries, we find (abbreviating $V_{QE\to Q'E'}$ as V) that

$$\begin{aligned} &F\left(\rho_{RQE}, \mathcal{R}_{Q\to QE}\left(\rho_{RQ}\right)\right) \\ &= F\left(V\rho_{RQE}V^\dagger, V\mathcal{R}_{Q\to QE}\left(\rho_{RQ}\right)V^\dagger\right) &(102)\\ &= F\left(\sigma_{RQ'E'}, V\mathcal{R}_{Q\to QE}\left(\rho_{RQ}\right)V^\dagger\right) &(103)\\ &\leq F\left(\sigma_{RQ'}, \mathrm{Tr}_{E'}\left\{V\mathcal{R}_{Q\to QE}\left(\rho_{RQ}\right)V^\dagger\right\}\right), &(104)\end{aligned}$$

where the inequality follows from monotonicity of fidelity under the discarding of subsystems. By defining the channel

$$\mathcal{E}_{Q\to Q'}(\cdot) \equiv \mathrm{Tr}_{E'}\left\{V_{QE\to Q'E'}\mathcal{R}_{Q\to QE}(\cdot)V^\dagger_{QE\to Q'E'}\right\}, \tag{105}$$

we find that

$$\varepsilon \geq -\log F\left(\sigma_{RQ'}, \mathcal{E}_{Q\to Q'}\left(\rho_{RQ}\right)\right), \tag{106}$$

establishing (98). ■

The following theorem provides a converse.

Theorem 8 *Suppose that the reduced dynamics are approximately CPTP, i.e., that to every unitary interaction $V_{QE\to Q'E'}$ leading to*

$$\sigma_{RQ'} = \mathrm{Tr}_{E'}\left\{V_{QE\to Q'E'}\rho_{RQE}V^\dagger_{QE\to Q'E'}\right\}, \tag{107}$$

there exists a CPTP map $\mathcal{E}_{Q\to Q'}$ such that

$$\frac{1}{2}\left\|\sigma_{RQ'} - \mathcal{E}_{Q\to Q'}\left(\rho_{RQ}\right)\right\|_1 \leq \varepsilon, \tag{108}$$

where $\varepsilon \in [0,1]$. Then the quantum data processing inequality is satisfied approximately, in the sense that

$$\begin{aligned} I(R;Q')_\sigma &\leq I(R;Q)_\rho \\ &\quad + 2\varepsilon\log|R| + (1+\varepsilon)\,h_2(\varepsilon/(1+\varepsilon)), \end{aligned} \tag{109}$$

and the conditional mutual information is nearly equal to zero as well:

$$I(R;E|Q)_\rho \leq 2\varepsilon\log|R| + (1+\varepsilon)\,h_2(\varepsilon/(1+\varepsilon)). \tag{110}$$

Proof. This follows directly from the assumption in (108), the Alicki–Fannes–Winter inequality, and the quantum data processing inequality:

$$\begin{aligned} &I(R;Q')_\sigma \\ &= H(R)_\sigma - H(R|Q')_\sigma &(111)\\ &= H(R)_{\mathcal{E}(\rho)} - H(R|Q')_\sigma &(112)\\ &\leq H(R)_{\mathcal{E}(\rho)} - H(R|Q')_{\mathcal{E}(\rho)} \\ &\qquad + 2\varepsilon\log|R| + (1+\varepsilon)\,h_2(\varepsilon/(1+\varepsilon)) &(113)\\ &= I(R;Q')_{\mathcal{E}(\rho)} \\ &\qquad + 2\varepsilon\log|R| + (1+\varepsilon)\,h_2(\varepsilon/(1+\varepsilon)) &(114)\\ &\leq I(R;Q)_\rho + 2\varepsilon\log|R| + (1+\varepsilon)\,h_2(\varepsilon/(1+\varepsilon)). &(115)\end{aligned}$$

The inequality for conditional mutual information follows the same reasoning we used to arrive at (100). ■

VII. CONCLUSION

We have shown how recent results regarding recoverability give enhancements to several entropy inequalities, having to do with entropy gain, information gain, disturbance, and complete positivity of open quantum systems dynamics. Our first result is a remainder term for the entropy gain of a quantum channel, which for unital channels is stronger than that which is obtained by directly applying the results of [30, 40]. This result implies that a small increase in entropy under a subunital channel is a witness to the fact that the channel's adjoint can be used as a recovery channel to undo the action of the original channel. Our second result regards the information gain of a quantum measurement, both without and with quantum side information. We find here that a small information gain implies that it is possible to undo the action of the original measurement (if it is efficient). The result also has operational ramifications for the information-theoretic tasks known as measurement compression without and with quantum side information. Our third result provides an information-theoretic measure of disturbance, introduced in [12], a strong operational meaning. We finally provide a robust extension of the main result of [9], establishing that the reduced dynamics of a system-environment interaction are approximately CPTP if and only if the data processing inequality holds approximately.

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[53] Notice that a trace-preserving positive map can be subunital only if the output space dimension is not smaller than the input space dimension, i.e., $\dim \mathcal{H}' \geq \dim \mathcal{H}$. In particular, if $\dim \mathcal{H}' = \dim \mathcal{H}$, then subunitality is equivalent to unitality.