Practical and reliable error bars for quantum process tomography

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(Received 20 August 2018; revised manuscript received 4 February 2019; published 8 May 2019)

Current techniques in quantum process tomography typically return a single point estimate of an unknown process based on a finite albeit large amount of measurement data. Due to statistical fluctuations, however, other processes close to the point estimate can also produce the observed data with near certainty. Unless appropriate error bars can be constructed, the point estimate does not carry any sound operational interpretation. Here, we provide a solution to this problem by constructing a confidence region estimator for quantum processes. Our method enables reliable estimation of essentially any figure of merit for quantum processes on few qubits, including the diamond distance to a specific noise model, the entanglement fidelity, and the worst-case entanglement fidelity, by identifying error regions which contain the true state with high probability. We also provide a software package, QPtomographer, implementing our estimator for the diamond norm and the worst-case entanglement fidelity. We illustrate its usage and performance with several simulated examples. Our tools can be used to reliably certify the performance of, e.g., error correction codes, implementations of unitary gates, or more generally any noise process affecting a quantum system.

DOI: 10.1103/PhysRevA.99.052311

I. INTRODUCTION

Quantum technologies are improving at an ever faster pace, not only by a concentrated academic effort, but increasingly via collaborations with industry. Quantum technologies require very precise manipulation and control of quantum systems, fueling the development of theoretical tools for precise calibration and characterization of quantum devices [1]. Notably, quantum-state tomography and quantum-process tomography (also known as quantum-process tomography) can infer the quantum state or the quantum process that describes a quantum device, providing a natural “quantum debugger” [2].

Quantum-state tomography aims to reconstruct the unknown state of a system with reference to a set of calibrated measurement apparatuses. Because of the intrinsically probabilistic nature of the outcomes of quantum measurements, one runs the experiment many times independently, collecting data from each run, and uses a statistical procedure to infer the quantum state. Formally, the reconstructed state is given by an estimator that associates a quantum state to a given data set. Statistical properties of estimators have been extensively studied not only in general for any statistical models [3–5], but also for quantum-state estimation [1,2,6]. A natural and popular estimator, which is furthermore computationally efficient, is the maximum likelihood estimator [7–10]. Error bars that are asymptotically optimal can be inferred from the quantum Fisher information [11], either via Cramér-Rao bounds or directly as confidence intervals [1,2,12–14].

In quantum-state estimation, one cannot always rely on results that are valid in the asymptotic regime of many measurements, and finite-size effects become important especially in settings where collecting measurements is costly [15–17]. Bayesian approaches avoid this problem by constructing credible regions that are well defined for any finite number of repetitions of the experiment [18,19]. An alternative approach that allows to make statements that are not prior dependent is to carefully and rigorously construct confidence regions with a precise confidence guarantee for finite number of copies. Such constructions have been demonstrated using regions based on a likelihood ratio [20], as well as procedures inspired by Bayesian methods in which a credible region can be “upgraded” to a confidence region [21,22].

Many tools for quantum-process tomography are adapted from quantum-state tomography, for instance, via the Choi-Jamiołkowski state process correspondence [23]. Beyond traditional process tomography [24], more advanced tools such as randomized benchmarking [25–28], gate-set tomography [29], and compressed sensing [30] display certain advantages such as a reduced number of required measurements. In the case of region estimators as in Refs. [21,22], some subtleties prevent a straightforward application of the corresponding tools for quantum states to quantum processes. Indeed, the set of quantum process is in one-to-one correspondence with only a subset of all bipartite states, namely, those whose reduced state on one system is maximally mixed; this constraint has to be incorporated explicitly in the region estimator. In this paper, we enrich the statistical toolbox for quantum-process tomography by providing a confidence region estimator for quantum process inspired by the state tomography method of Christandl and Renner [21], with a rigorous confidence guarantee for any finite number of measurements.

Often in certifying specific applications, we are not interested in the full knowledge of the quantum process; a property of the unknown channel suffices. For example, in quantum key distribution we are often interested in how close the final
state output by the protocol is to the ideal key state; this is captured for instance by the fidelity or the trace distance of the real state to the ideal state [31]. Likewise, in quantum computing a relevant figure of merit that enables fault-tolerant computation is the error threshold captured by the diamond distance or the worst-case entanglement fidelity of the real implemented gate relative to the ideal gate [32]. Note that a bound on the diamond distance or the entanglement fidelity to a given fixed channel may confine the true channel to a small region in channel space.\footnote{If after tomography one obtains the statement |Λ − Λ^{\text{ideal}}|_\diamond ≤ 10^{-3} with 99\% confidence, then for all practical purposes we can consider that Λ = Λ^{\text{ideal}}. Thus, knowledge of the figure of merit may imply knowledge of all the entries in the process matrix of the unknown channel Λ.} For these reasons, and because this significantly simplifies our analysis, we focus on estimators for quantum processes that report confidence intervals for a given figure of merit.

Summary of main results. Our main contribution is three-fold:

(i) A confidence region estimator for channel tomography through the use of the Christandl-Renner-Faist estimator for states and the Choi-Jamiolkowski isomorphism between quantum states and quantum processes. We call this the bipartite-state sampling method.

(ii) A confidence region estimator to directly (without first tomographing the Choi state associated with the channel) estimate quantum processes and its proof of correctness. We call this the channel-space sampling method.

(iii) A software package called QTomographer [33] accompanying our theoretical results for analyzing experimental data. Our software returns quantum error bars which capture all the information about the unknown channel derivable from the tomographic data and enable the user to construct confidence regions for any confidence level of interest.

By comparing the differences of the two estimators, we obtain a better understanding about the relationship between probability measures on state space and channel space which may be of independent interest. Because the estimators return a confidence region, they will work without any assumption on the prior distribution of the unknown process.

To illustrate how to use our result, we consider the scenario of certifying a quantum memory (an example of quantum property testing [34]). This corresponds to certifying that a quantum device (approximately) implements the identity channel. We consider three possible figures of merit: the diamond distance to the identity channel, the entanglement fidelity, and the worst-case entanglement fidelity [35,36]. Our method yields a reliable estimation of these figures of merit.

The paper is organized as follows. We first demonstrate in Sec. II how one can use our method to obtain reliable information in a tomography experiment. The correctness of our tools is justified in Sec. III where we present the main results. Then, we study the behavior of our numerical implementations in Sec. IV before concluding our paper with future directions (Sec. V). We leave the formal statements and detailed derivations of our results to the Appendices.

![FIG. 1. The workflow of rigorous process tomography. Our data analysis QTomographer supports both prepare-and-measure and ancilla-assisted experimental schemes. The conclusion is guaranteed without any prior information on the unknown quantum process.](image)

![FIG. 2. Illustrations of (a) prepare-and-measure and (b) ancilla-assisted tomographic schemes for an unknown channel Λ_{A→B}. In prepare and measure, one can only prepare input state $\sigma_A^j$ which is fed into an unknown channel whose output state is measured by some POVM with elements $E_B$. In ancilla assisted, one can prepare entangled input state with some reference system $P$, and measure jointly the output using some POVM with elements $E_{B'P'}$.](image)
The quantum memory, and enable us to construct confidence intervals skewness. These data encode information about the performance of peak, $\Delta_1$ which we called $\Delta$, which is compactly described by the triple of numbers $(v_0, \Delta, \gamma)$.

The quantum error bars are perfectly analogous to classical error bars: the latter are indeed a concise, intuitive description of the error from which one easily determines rigorous confidence intervals. For this reason it is a natural object to report at the end of a process tomography procedure.

If one wishes to actually derive rigorous confidence regions for the diamond norm distance, one may proceed as follows. First, one fixes a confidence level, say $\alpha = 99\%$, which sets the corresponding error parameter as $\epsilon = 1 - \alpha = 10^{-2}$. By Theorem 2, for $n = 45000$ (size of our data set $E$) and $d_A = d_B = 2$, we need to find a region of diamond norm distance values with weight at least

$$1 - \frac{\epsilon}{2} \left( \frac{2n + d_A^2 d_B^2 - 1}{d_A^2 d_B^2 - 1} \right)^{-2} \geq 1 - 10^{-151}.$$ 

With reference to Fig. 3, this means we need to find the position such that the area under the curve exceeds $1 - 10^{-151}$. A numerical integration leads to a region at least as large as $[0, 0.24]$. Together with the enlargement by

$$\delta = \sqrt{\frac{2}{n} \left( \ln \frac{2}{\epsilon} + 3 \ln \left( \frac{2n + d_A^2 d_B^2 - 1}{d_A^2 d_B^2 - 1} \right) \right)} = 0.1$$

(to exclude nearby channels which could result in the same observed data set with high probability) the final confidence region is $[0, 0.34]$. This means we have certified that the diamond norm distance of the unknown quantum memory to an ideal quantum memory is at most $0.34$ with $99\%$ confidence. In general, increasing the number $n$ of measurement data points will shrink this confidence interval [due to the exponential decays in the diamond distance density (see also Appendix F)].

We emphasize that the unnaturally large size of the regions is due in large part to a technical difficulty in the proofs of our bounds that is dealt with by employing tools that are known not to be tight in this context. For this reason, the quantum error bars are more informative than the actual final confidence regions.

This concludes the general workflow associated with our tomographic tools. The next section explains at a high level how our software transform tomographic data into confidence regions.

**III. MAIN RESULTS**

Our software package QPtomographer is built on top of two rigorously proven theoretical constructions. These are confidence region estimators based on the bipartite-state sampling method or the channel-space sampling method. The bipartite-state method works in the ancilla-assisted scheme, while the channel-space method works in both ancilla-assisted and prepare-and-measure schemes. This section gives a high-level overview of the constructions together with the main

![Figure 3](image-url)
ideas behind the proof of correctness, and leave the details to Appendices B and C, respectively. We begin with a brief motivation for confidence region estimators.

A. Confidence region estimators

In the limit of infinite data (i.e., the number of records in data set $E$ is infinity), it is possible to exactly compute the probabilities of each measurement outcome from $E$ and reconstruct the unknown channel by linear inversion on these observed probabilities [2]. However, in the practical scenario of finite data (i.e., data set $E$ contains $n$ records) statistical fluctuations will imply the failure of all point estimation methods such as linear inversion or maximum likelihood estimation. This is due to the fact that channels close to the point estimate can produce the same data set with high probability.

In order to make statistically rigorous and operationally sound statements on the unknown channel in this regime, we turn to region estimators, which are generalizations of the process of constructing error bars. We will look at a type of region estimators known as confidence region estimators. These are maps from data $E$ to subsets $S_E \subseteq \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B)$ of the set of quantum processes with the property that for all $\Lambda \in \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B)$

$$
\Pr_E[\Lambda \in S_E] \geq \alpha,
$$

where $\alpha$ is a prefixed confidence level and the probability is evaluated over the random data set $E$ according to the distribution $\Pr[\mathcal{E}|\Lambda]$. It is important to note that confidence is a property of the entire estimator (the procedure $E \mapsto S_E$) and not of any particular subset $S_E$ produced by the estimator.

The operational meaning of confidence region estimators can be understood as follows. Suppose the black box implementing the unknown channel $\Lambda_{\text{true}}$ is in fact prepared by a referee, who knows exactly to which channel the black box applies. We proceed with a sequence of state preparations, applications of the channel, and measurements of the output states to obtain a data set $E$. Then, we apply the estimator on $E$ to get $S_E$. Repeating this procedure a large number of times, say $N = 10^3$, if $r$ denotes different repetitions, then we obtain different data sets $E(r = 1), \ldots, E(r = 10^3)$ with corresponding conclusions that the true channel $\Lambda_{\text{true}}$ should be in the region $S_{E(r=1)}, \ldots, S_{E(r=10^3)}$. Now, since the referee knows exactly the unknown channel, the referee can evaluate the proportion of correct conclusions

$$
\frac{|\{r : r = 1, \ldots, N \text{ and } \Lambda \in S_{E(r)} \text{ is true}\}|}{N}.
$$

If the estimator used is a confidence region estimator with confidence level $\alpha = 0.99$, then in the limit of $N \rightarrow \infty$ this proportion is at least 0.99. This is the meaning of confidence: the correct conclusion is guaranteed for a large number of uses of the estimator, regardless of the unknown channel. Note that for a specific use of the estimator which returns $S_E$, we cannot draw the conclusion that $\Lambda \in S_E$.

An alternative justification of confidence regions comes from a Bayesian point of view: Bayesian tomography uses outcomes of measurements to update a prior distribution about the quantum state to a posterior distribution. While this posterior clearly depends on the prior, it is known that when enough data are collected, the posterior distribution is no longer sensitive to the exact prior which was originally used (as long as the original prior has full support). Now consider a high-weight region of a posterior distribution, which is also known as a credible region. We may ask to what extent this region remains a credible region if we change the underlying prior. It turns out that for a large enough number of measurements, we may find regions which are credible regions for any prior, except for some exceptionally unlikely measurement data sets [21]. Such regions are precisely confidence regions.

B. Our confidence region estimators

Our method of constructing region estimators uses the information about the underlying unknown channel via the likelihood function defined generically for an observed data set $E$ as

$$
\mathcal{L}(\Lambda|E) = \Pr(E|\Lambda),
$$

where the probability of the data set $E$ under the assumption that the unknown channel is $\Lambda$ is given by Born’s rule. The specific form of the likelihood function depends on the scenarios and assumptions we postulate (cf. Appendices B and C). The likelihood function can be seen as giving a ranking about which channel best produce the observed data set. We now present our methods of process tomography.

Bipartite-state sampling method. The main idea behind this method is that a quantum process is in correspondence with bipartite Choi states via the Choi-Jamiolkowski isomorphism. Hence, we can construct confidence regions for quantum states using the method of Christandl-Renner, and then perform an additional classical post-processing step to recover a confidence region for quantum processes.

Let us now first assume the use of an ancilla-assisted tomographic scheme, Fig. 2(b), which loosely corresponds to physically performing the Choi-Jamiolkowski isomorphism in the laboratory. This means having access to a full-rank bipartite-entangled state $|\psi_{AP}\rangle$ as input to the channel, and performing tomography on the output state $\rho_{BP} := \Lambda_{A-B}(\psi_{AP})$ which is the unknown Choi state associated with the unknown channel.

Treating $\rho_{BP}$ as the unknown state in a state tomography problem, we now apply the Christandl-Renner method of constructing confidence regions from tomographic data. Recall that the Christandl-Renner confidence region is constructed from the measure

$$
d\mu_E(\sigma_{AB}) := c_E^{-1} \text{tr}(\sigma_{AB}^{\otimes n} E) d\sigma_{AB},
$$

where $c_E = \int \text{tr}(\sigma_{AB}^{\otimes n} E) d\sigma_{AB}$ is the normalizing constant, and $d\sigma_{AB}$ is the uniform distribution on bipartite density matrices (obtained by tracing out a Haar random pure state on a larger space). Note that $\text{tr}(\sigma_{AB}^{\otimes n} E)$ is the likelihood function for the outcome $E$ given the state $\sigma_{AB}$ in this scenario. Confidence regions for the unknown $\rho_{AB}$ can be constructed from $d\mu_E(\sigma_{AB})$ as the following proposition asserts.

Theorem 1 (Christandl and Renner [21], informal). Let $n$ be the number of systems measured by a POVM during tomography and $1 - \epsilon$ be the desired confidence level. Let $S_{\mu_E} \subseteq D(\mathcal{H}_{AB})$ be any set of bipartite states with high weight under the probability measure $d\mu_E(\sigma_{AB})$. Then, the
enlargement in purified distance $S_{\mu e}^d$ where

$$\delta = \sqrt{\frac{2}{n} \left( \ln \frac{2}{\epsilon} + 2 \ln s_{2n,d}^2 \right)}$$

with $s_{n,d} := (n^{d-1})$ is a confidence region of confidence level $1 - \epsilon$.

Intuitively, we can think of the enlargement as a way to exclude nearby states or channels (relative to a proposed region of states or channels) that can give rise to the same observed data set $E$ with nonzero probability.

The confidence region $S_{\mu e}^d$ contains bipartite quantum states which are not Choi states. This is due to the fact that the method of Christandl and Renner does not a priori allow the Choi state constraint $\text{tr}_B(\sigma_{AB}) = 1\Lambda/dA$. Hence, we have to invent an additional post-processing step to map $S_{\mu e}^d$ to a region consisting of exclusively Choi states. By the Choi-Jamiolkowski isomorphism we then have a confidence region for the unknown quantum process. The detailed explanation is left to Appendix B.

Channel-space sampling method. This method is a construction of confidence region that directly returns channel-space confidence regions. Compared to the bipartite-state method, the channel-space method works in both the prepare-and-measure and ancilla-assisted tomographic schemes and takes into account the a priori knowledge that we are estimating a quantum process. This leads to computational efficiency relative to the bipartite-state method because the additional post-processing step of the bipartite-state method is not required here.

The estimator is constructed from the probability measure on the set of quantum processes $E(\mathcal{H}_A \to \mathcal{H}_B)$:

$$d\nu_E(\Lambda) := c_E^{\epsilon} \mathcal{L}(\Lambda | E) d\nu(\Lambda),$$

where $\mathcal{L}(\Lambda | E)$ is the likelihood for the event $E$ given a channel $\Lambda$, $c_E = f \mathcal{L}(\Lambda | E) d\nu(\Lambda)$ serves as a normalizing constant, and $d\nu(\Lambda)$ is the Haar-induced measure on $E(\mathcal{H}_A \to \mathcal{H}_B)$. The likelihood function is adapted depending on prepare-and-measure or ancilla-assisted tomographic scheme and is defined as the probability of obtaining the data set $E$ given a channel $\Lambda$. Informally, this measure captures the information of the unknown channel as revealed by the observed data set $E$ in an unbiased manner (that is, without using any prior knowledge on the unknown).

Given this measure, we obtain the following:

Theorem 2 (informal). Let $n$ be the number of channel uses during tomography and $1 - \epsilon$ be the desired confidence level. Let $R_{\mu e} \subseteq E(\mathcal{H}_A \to \mathcal{H}_B)$ be a set of channels with high weight under the probability measure $d\nu_E(\Lambda)$. Then, the enlargement in purified distance (for quantum processes, induced from states) $R_{\mu e}^d$ where

$$\delta = \sqrt{\frac{2}{n} \left( \ln \frac{2}{\epsilon} + 3 \ln s_{2n,d}^2 \right)}$$

with $s_{n,d} := (n^{d-1})$ is a confidence region with confidence level $1 - \epsilon$.

Confidence interval for figures of merit. In practice, we choose the region in Theorem 2 for any chosen figure of merit to be the subset of channels whose figure of merit is better than a certain threshold. For the diamond norm distance to the ideal channel, we consider

$$R = \{ \Lambda : |\Lambda - \Lambda_{\text{ideal}}|_{\diamond} \leq \gamma_E \},$$

and for the worst-case entanglement fidelity we consider

$$R = \{ \Lambda : F_{\text{worst}}(\Lambda) \geq \gamma_E \}.$$
The output of our classical data analysis is called “quantum error bars” which contain all the information about the figure of merit that can be obtained from the tomographic data set. From here, it is easy to construct confidence regions for any specified confidence level.

**A. One-qubit example**

In the following, see example qubit-noisy-identity in our software package [33] for more details.

**Step 1. Data collection:** Consider the scenario of testing the performance of a quantum memory \( \Lambda_{A \to B} \). The ideal channel we wish to implement is the identity channel \( I \). Suppose that the real channel implemented in the experiment is the depolarizing channel

\[
\Lambda_{A \to B}(\rho) = p \rho + (1 - p) d_B^{-1}B,
\]

acting on one qubit \( (d_A = d_B = 2) \), with the parameter \( p = 0.9 \). In other words, the experiment is slightly off from the ideal implementation by some white noise.

Furthermore, we consider the ancilla-assisted scheme, and assume that the input to the channel is half of a pure entangled state \( |\psi\rangle_A = (\sigma_A^{1/2} \otimes I) d_A^{1/2} |\Phi\rangle_A \), where we choose

\[
\sigma_A = \begin{pmatrix} 0.6 & 0.1 \\ 0.1 & 0.4 \end{pmatrix},
\]

which mimics an input state which deviates slightly from the maximally mixed state. Note that the entangled input state has full Schmidt rank.

Since we do not have an actual experiment, we have to simulate Pauli measurements on the joint state \( \rho_{BP} \) after application of the channel \( \Lambda_{A \to B} \), with two possible outcomes for each of the three measurement settings. For each measurement setting, 500 measurement outcomes were simulated. These constitute the information contained in the (simulated) observed data set \( E \) with \( n = 45 \, 000 \).
We now subject these data set to an analysis which we aim to measure three figures of merit corresponding to our unknown channel: the diamond distance to the identity channel, the average entanglement fidelity, and the worst-case entanglement fidelity. Refer to the Appendix A for the precise definitions.

Steps 2 and 3. Random sampling and histogram: We run the channel-space sampling method to estimate the three figures of merit. The calculation of all three figures of merit functions was done in C++ using the SCS toolbox \[39,40\]. A simple PYTHON interface was used to control the execution of the program. The Metropolis-Hastings random walk is run on the space of all quantum processes using elementary rotation jump distribution, as described in Appendix D, until 32 768 data points have been collected. Samples from the random walk allow to construct a histogram approximating a specific distribution of the figure of merit.

Step 4. Fit analysis of histograms: In our example, we discovered that the fit model as described in (9) does have good agreement with the underlying histogram bins, as underscored by goodness-of-fit values (reduced $\chi^2$) of the order of $\sim 4$ in most cases. In case of poor fit to the default model, one may modify the empirical fit model until an acceptable goodness of fit is achieved, before quantum error bars and confidence regions can be derived.

Step 5. Quantum error bars and confidence regions: The quantum error bars $(\nu_0, \Delta, \gamma)$ are a simple translation from the parameters of the fit model. The steps toward a confidence region for diamond norm have been illustrated in Sec. II. In theory, we have the guarantee that collecting a larger data set will yield smaller regions converging to the true value. Unfortunately, the confidence interval for diamond norm distance returned by our method is unreasonably large for the current example: for 99% confidence level we are able to bound the diamond norm by 0.34 as compared to the true value of 0.05. We believe that this is due to operator inequality involved in bounding the failure probability (Proposition 1). Further research is needed to provide better construction of confidence regions (i.e., more efficient in terms of the number of data samples $n$).

B. Two-qubit example

Now, we consider a two-qubit example to illustrate the practicality of our method in this situation. This example also shows that the channel-space and the bipartite-state sampling methods do not in general produce the same histogram, as one might have concluded from Fig. 4.

In the following, see example 2qubits-noisy-identity in our software package [33] for more details. Suppose that the real channel implemented in the experiment the two-qubit depolarizing channel

$$\Lambda_{A\rightarrow B}(\rho) = p \rho + (1-p) d_B^{-1/2} B,$$

with $d_A = d_B = 4$, and we are interested in the diamond distance to the identity channel. Assuming access to state preparation that produces $|\psi\rangle_{AP} = (\sigma_A^{1/2} \otimes 1) d_A^{1/2} |\Phi\rangle_{AP}$ with

$$\sigma_A = \begin{pmatrix} 0.35 & 0 & 0.04 & 0.1i \\ 0 & 0.15 & 0.05 & 0 \\ 0.04 & 0.05 & 0.32 & 0 \\ -0.1i & 0 & 0.18 & 0 \end{pmatrix},$$

and $3^4 = 81$ Pauli measurement settings each having $2^2 = 4$ outcomes. We perform similar analyses on a simulated data

![Distribution of figures of merit for the single-qubit noisy quantum memory example](image-url)
set of size $n = 40,500$ which we generated using the state preparations and measurements described above. The result is presented in Fig. 5.

The channel-space sampling method’s $h(v)$ is peaked at lower values of the figure of merit, as can be seen in Fig. 5. We observe that, in this case, the knowledge of the input state significantly shifts the corresponding histogram distribution toward lower values of the figure of merit, allowing to construct smaller confidence regions. Based on several examples studied, this is not always the case; with less noise (smaller $p$), for instance, the curve for $\mu(v)$ and the curve for $h(v)$ get closer to each other.

On a technical level, we show that the Hilbert-Schmidt measure over the bipartite states factorizes as a measure over states on the input system and the relevant measure over all channels (Appendix G). Hence, a large uncertainty over the input state may enlarge the resulting region as opposed to considering a region only on the channel space for a fixed known input state. However, it is not impossible that under some lucky circumstances a finite distribution width on the input state helps add more weight to regions of a higher figure of merit, effectively shrinking the region. Indeed, it could happen that the input state assumed in the channel-space method is far from the optimal state for distinguishing the channels in terms of the diamond norm; in such a case a prior which is more “smeared out” over different input states might result in smaller quantum error bars for the diamond norm. We believe that this is why neither method performs globally better than the other. See Appendix G for further details on the relationship between the two methods.

V. CONCLUSIONS

One might think that carrying over the notion of quantum error bars in quantum-state tomography to quantum-process tomography is as straightforward as converting quantum states to channels via the Choi-Jamiołkowski isomorphism.

However, our study reveals a more complicated structure. We find that different analysis methods are suited to different experimental quantum-process tomography setups. In the experimentally more realistic prepare-and-measure scheme, a judicious use of the prior knowledge about the input state to the process allows us in typical situations to obtain tighter quantum error bars for the process. These results are obtained by developing a method, along with corresponding proofs, which are specific to process tomography.

On the other hand, in the case of the ancilla-assisted scheme, we can directly apply the methods developed for quantum-state tomography, harnessing them to directly yield reliable statements about the quantum process itself, while ignoring any information the measurements provide about the input state used to probe the process.

We hence provide a fully fledged and practical toolbox named \textit{QPtomographyer}, with solid theoretical foundations, for quantum-process tomography of arbitrary quantum processes, using any experimental quantum-process tomography setup, and given measurement outcomes from any measurement settings. Our software package facilitates the numerical analysis in practice by automating the implementation of the Metropolis-Hastings random walk, as well as the calculation of the diamond norm, by simple high-level PYTHON function calls, while transparently delegating the computation-intensive routines to heavily optimized C++ code which makes use of modern programming techniques including template metaprogramming and exploiting hardware SIMD instructions.

On the spectrum of characterization tools for quantum devices, our method can be seen as lying on the opposite end of randomized benchmarking \cite{25–28}. While slightly more involved, our technique can be applied to any choice of state preparations and measurements, and can be applied to any individual process. By determining the diamond norm or the worst-case entanglement fidelity to any given ideal process, we provide individual full characterization of the processes implemented by individual gates. More generally, our methods allow the reliable estimation of any specific property of the quantum process.

We note that our method is currently limited to processes acting on few qubits, as our confidence region produces unreasonably large regions, and the algorithm stores dense representations of the quantum process. However, we expect that our methods will be used to certify individual components of complex setups, for instance, individual two-qubit gates. Because we estimate robust, composable figures of merit such as the worst-case entanglement fidelity or the diamond norm, the composition of individually certified components is still certified to function accurately. The generous size of the region is partially due to the fact that we provide rigorous confidence guarantees that are valid for any finite number of measurements, the proof of which requires the use of mathematical bounds that are not necessarily tight. In
practice, however, it appears that the quantum error bars are comparable in magnitude to error bars obtained by methods that would formally be exact only in the asymptotic regime, such as the bootstrap [22]. As future work we would envisage a more systematic analysis of the discrepancies between the quantum error bars in the finite regime and traditional methods such as those based on the Fisher information or the bootstrap.

In addition, we may ask whether the channel method is always superior to the bipartite sampling method. As noted above, the additional prior knowledge about the input state which the channel-space method enjoys in contrast to the bipartite sampling method is not sufficient to guarantee this. We leave a more precise understanding of the relation between our two methods open for future study.

ACKNOWLEDGMENTS

L.P.T., J.H., D.E., and S.W. are supported by an ERC Starting Grant (S.W.), an NWO VIDI Grant (S.W.), and an NWO ZwaarteKraft grant (Q.S.C.). P.F. acknowledges support from the Swiss National Science Foundation (SNSF) through the Early PostDoc. Mobility Fellowship No. P2EZP2_165239 hosted by the Institute for Quantum Information and Matter (IQIM) at Caltech, from the IQIM which is a National Science Foundation (NSF) Physics Frontiers Center (NSF Grant No. PHY-1733907), and from the Department of Energy Award No. DE-SC0018407.

APPENDIX A: NOTATIONS AND PRELIMINARIES

We begin by setting up some notations and recalling standard definitions. For more information on states and processes see [41–43].

1. Quantum processes and figures of merit

Let \( \mathcal{H}_A \) be the Hilbert space of dimension \( d_A \) associated with the quantum system denoted \( A \). By \( \text{D}(\mathcal{H}_A) \) we mean the subset of \( \text{End}(\mathcal{H}_A) \), the set of linear transformations on \( \mathcal{H}_A \), consisting of density matrices \( \rho_A \geq 0 \) (positive semidefinite) with \( \text{tr}(\rho_A) = 1 \). Composite systems are described by tensor product constructions, for instance, \( \mathcal{H}_{AB} = \mathcal{H}_A \otimes \mathcal{H}_B \) is the Hilbert space of composite system \( AB \).

Quantum measurements on quantum system are the positive operator valued measures or POVMs on \( \mathcal{H} \). For finite number of outcomes, a POVM is a set of positive operators, the effects, that sum to the identity operator on \( \mathcal{H} \). We will overload the notation \( E \) to mean an outcome label, and also the effect \( E \) (i.e., an operator or matrix) in the POVM. This is equivalent to the usual “observables” formulation of measurement, i.e., a Hermitian operator. For example, a \( Z \) measurement or observable has two outcomes \( E = +1 \) and \( E = -1 \) with associated effects \( |0\rangle \langle 0| \) and \( |1\rangle \langle 1| \), respectively.

A quantum process \( \Lambda_{A\rightarrow B} \) mapping a quantum system \( A \) to a quantum system \( B \) is a completely positive trace-preserving linear map from \( \text{End}(\mathcal{H}_A) \) to \( \text{End}(\mathcal{H}_B) \). In general, we will denote quantum processes by capital Greek letters. We will often drop the subscripts when the quantum systems are clear from the context.

The set of all possible quantum processes is denoted \( \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B) \), and it is in one-to-one correspondence with the set of bipartite Choi states \( \mathcal{C}(\mathcal{H}_{AB}) \) via the Choi-Jamiolkowski isomorphism

\[
J : \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B) \rightarrow \text{End}(\mathcal{H}_A \otimes \mathcal{H}_B),
\]

\[
\Lambda_{A\rightarrow B} \mapsto (I_A \otimes \Lambda_{A\rightarrow B})(|\Phi\rangle \langle \Phi|_A).
\]

Throughout the Appendices, we will use the convention that \( \Lambda_{AB} \) is the Choi state associated with the channel \( \Lambda_{A\rightarrow B} \).

The action of the channel can be recovered from its Choi state by the inverse of Choi-Jamiolkowski isomorphism \( \Lambda(\rho) = d_A \text{tr}_A(\Lambda_{AB} : \rho_B^{\top} \otimes 1_A) \).

Recall that the fidelity between two states \( \sigma, \sigma' \) is defined as

\[
F(\sigma, \sigma') := \| \sqrt{\sigma} \sqrt{\sigma'} \|_1 = \text{tr} \sqrt{\sqrt{\sigma} \sigma' \sqrt{\sigma}},
\]

and the purified distance between quantum states is defined as \( P(\sigma, \sigma') := \sqrt{1 - F(\sigma, \sigma')^2} \). Then, the purified distance between channels is defined as

\[
P(\Psi_{A\rightarrow B}, \Psi'_{A\rightarrow B}) := P(\Psi_B, \Psi'_B).
\]

a. Diamond distance

We first introduce the familiar diamond distance. The diamond distance from channel \( \Lambda_{A\rightarrow B} \) to the identity channel of \( \mathcal{H}_A \) is the maximal diamond distance of \( \Lambda_{A\rightarrow B} \) and \( \text{id}_A \). Explicitly, the set of Choi matrices is defined as \( \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B) \) under the Choi-Jamiolkowski isomorphism and has the following compact description:

\[
\mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B) = \{ \sigma : \sigma \in D(\mathcal{H}_A) \text{ and } \text{tr}(\sigma) = 1 \}.
\]

Throughout the Appendices, we will use the convention that \( \Lambda_{AB} \) is the Choi state associated with the channel \( \Lambda_{A\rightarrow B} \).

The action of the channel can be recovered from its Choi state by the inverse of Choi-Jamiolkowski isomorphism

\[
\Lambda(\rho) = d_A \text{tr}_A(\Lambda_{AB} : \rho_B^{\top} \otimes 1_A).
\]

The diamond distance between channels is defined as

\[
d_{\mathcal{C}}(\Lambda_A, \Lambda'_A) := \| \Lambda_A - \Lambda'_A \|_{\mathcal{C}}.
\]

This function \( d_{\mathcal{C}}(\Lambda_A, \Lambda'_A) \) can be cast as a semidefinite program [44]

\[
\text{Primal problem}
\]

maximize: [\( d_A \Lambda_{AB} - d_A \Lambda_{AB}^{\text{ideal}}, W \)]

subject to: \( W \leq 1_{\mathcal{B}} \otimes \rho \),

\( W \geq 0 \),

\( \rho \in \text{D}(\mathcal{X}) \).

\text{Dual problem}

minimize: \( \| \text{tr}_B(Z) \|_{\infty} \)

subject to: \( Z \geq d_A \Lambda_{AB} - d_A \Lambda_{AB}^{\text{ideal}} \),

\( Z \geq 0 \).

b. Entanglement fidelity

The entanglement fidelity is another measure of how close a given channel is to the identity channel. More specifically, it measures how well a channel preserves the maximally
entangled state. The entanglement fidelity of a channel \( \Lambda_{A \rightarrow B} \) with \( B \simeq A \) is defined as
\[
F_e(\Lambda) = F^2(\Lambda_{A \rightarrow B}(\Phi_{AA}), \check{\Phi}_{AB}), \tag{A7}
\]
recalling that \( |\Phi_{AA}\rangle \) is the normalized maximally entangled state between the systems \( A \) and \( \check{A} \).

Because \( \Lambda_{AB} = \Lambda_{A \rightarrow B}(\Phi_{AA}) \) is the normalized Choi state corresponding to the channel \( \Lambda_{A \rightarrow B} \), the entanglement fidelity of the channel \( \Lambda_{A \rightarrow B} \) is in fact exactly the fidelity of the corresponding normalized Choi state to the maximally entangled state:
\[
F_e(\Lambda) = F^2(\Lambda_{AB}, \check{\Phi}_{AB}). \tag{A8}
\]

\textbf{c. Worst-case entanglement fidelity}

The worst-case entanglement fidelity is a better measure of the reliability of the channel to simulate the identity channel, if we have to worry about any possible input state being fed into the channel. In effect, the worst-case entanglement fidelity measures how well the channel preserves any given state on a system and any purification. It is defined as
\[
F_{\text{worst}}(\Lambda_{A \rightarrow B}) = \inf_{\Phi_{AB}} F^2(\Lambda_{A \rightarrow B}(\Phi_{AA}), \sigma_{AB}), \tag{A9}
\]
where the optimization ranges over all bipartite quantum states \( \sigma_{AB} \) defined over the input \( \check{A} \) and a reference system \( A \simeq \check{A} \). The optimization variable, which appears in both slots of the fidelity \( F \), may be restricted to pure states without loss of generality.

Now, we show that the worst-case entanglement fidelity can be computed by evaluating a simple semidefinite program. That a semidefinite program formulation of the worst-case entanglement fidelity can be used in the context of quantum error correction to find suitable recovery procedures for fixed input was put forth in Refs. [45,46]. We build upon those constructions to optimize over the input state, while in our case the problem is simplified as there is no recovery operation. Using our notation, we write
\[
F_{\text{worst}}(\Lambda_{A \rightarrow B}) = \inf_{\Phi_{AB}} \frac{1}{\Phi_1} \text{tr}(\Lambda_{A \rightarrow B}(T_A \check{\Phi}_{AA} T_A^\dagger) T_A \check{\Phi}_{AB} T_A^\dagger), \tag{A10}
\]
where we have defined the non-normalized maximally entangled state \( |\Phi_{AB}\rangle = d_A^{1/2} |\check{\Phi}_{AB}\rangle \). The last equality comes from the fact that any bipartite pure state \( |\phi_{AA}\rangle \) can be parametrized by a complex matrix \( T_A \) satisfying \( \text{tr}(T_A T_A^\dagger) = 1 \) via \( |\phi_{AA}\rangle = T_A |\check{\Phi}_{AA}\rangle \), with moreover \( T_A^\dagger = T_A \). The second measure is defined on the set of quantum systems and any purification. It is defined as
\[
F_{\text{worst}}(\Lambda_{A \rightarrow B}) = \inf_{\Phi_{AB}} F^2(\Lambda_{A \rightarrow B}(\Phi_{AA}), \phi_{AB}), \tag{A11}
\]
This is a minimization over a positive-semidefinite quadratic form in \( \rho_A |\Phi_{AB}\rangle \), so it is (quite surprisingly) a convex optimization in terms of \( \rho_A \). We know that positive-semidefinite quadratic optimizations may be written as semidefinite programs. Indeed, for any positive-semidefinite matrix \( Q = MM^\dagger \), we have that \( \langle \psi | Q | \phi \rangle \leq \mu \) if and only if \( \begin{bmatrix} \mu \\ \frac{1}{\mu} - \frac{M}{\mu} \end{bmatrix} \geq 0 \). So, finally, we may write the worst-case entanglement fidelity as a semidefinite program in terms of the real variable \( \mu \) and the positive-semidefinite variable \( \rho_A \geq 0 \):
\[
\begin{bmatrix} \mu \\ \text{tr}(\rho_A) = 1 \\ (|\Phi_{AB}\rangle \rho_A M_{AB}) \begin{bmatrix} M_{AB}^\dagger \rho_A |\Phi_{AB}\rangle \\ \mu \end{bmatrix} \end{bmatrix} \succeq 0, \tag{A12}
\]
where \( D(\Lambda_{AB}) \) admits a purification space \( \text{Pure}(H_{ABA'B'}) \) whose elements are rank-one density operators on \( H_{ABA'B'} \) with \( A'B' \) being an isomorphic copy of \( AB \). The Haar measure \( dU_{ABA'B'} \) then induces a measure on \( \text{Pure}(H_{ABA'B'}) \) via the relation \( |\psi \rangle \langle \psi | = U |\psi_0\rangle \langle \psi_0 | U^\dagger \) for an arbitrary pure state \( |\psi_0\rangle \), which induces a measure \( d\sigma_{AB} \) on \( D(\Lambda_{AB}) \) by partial tracing. The second measure is defined on the set of quantum processes, or equivalently on the set of bipartite Choi states. Let
\[
\mathcal{P} \subseteq \left\{ |\psi\rangle \in H_{ABA'B'} : \text{tr}_{B'A'}(|\psi\rangle\langle \psi |) = d_A^{-1} |\check{A}\rangle \right\}, \tag{A14}
\]
be the set of purifications of arbitrary Choi states. Without loss of generality, let us define a fixed reference state in \( \mathcal{P} \),
\[
|\check{\sigma}_0\rangle := \frac{1}{d_A} \sum_{i=1} d_A |i\rangle A| v_i\rangle_{B'A'} \tag{A15}
\]
with \( \{|\psi\rangle_B\rangle\} \) some fixed orthonormal set of vectors. Then, for all \( |\psi\rangle \in \mathcal{P} \), there exists a unitary \( U_{B\rightarrow B} \) such that \( |\psi\rangle = \mathcal{I}_B \otimes U_{B\rightarrow B}(|\psi_0\rangle) \). This relation transfers the unique Haar measure \( dv(U_{B\rightarrow B}) \) on the unitary group \( U(\mathcal{H}_B) \) to a measure on \( \mathcal{P} \) which we will denote as \( dv(|\psi\rangle) \). Again, by partial tracing the system \( A' B' \), this measure induces the measure \( dv(|\psi_{AB}\rangle) \) on Choi states \( \mathcal{E}(\mathcal{H}_{AB}) \) [also denoted as \( dv(|\psi_{A\rightarrow B}\rangle) \) by changing the dummy variable]. Finally, taking the inverse of the Choi-Jamiolkowski isomorphism gives the induced measure \( dv(|\psi_{A\rightarrow B}\rangle) \) [or in a different notation \( dv(|\Lambda_{A\rightarrow B}\rangle) \)] on channel space \( \mathcal{E}(\mathcal{H}_A \rightarrow \mathcal{H}_B) \) which is the starting point of the channel-space sampling method.

The relation between these measures will be discussed in Appendix G when we compare the two region estimators.

### APPENDIX B: BIPARTITE-STATE SAMPLING METHOD

This method requires experimentalists to work in the ancilla-assisted scheme [see Fig. 2(b)]: we select a full Schmidt rank entangled state \( \rho_{AP} \), a collection of bipartite measurements \( E^{(i)}_k \) with corresponding effects \( E^{(i)}_k(\cdot) \), and assume the experiment can implement the channel \( \Lambda_{A' B'} \mathcal{I} \), where \( \mathcal{I} \) is the identity map. Again, the collection of measurements should be informationally complete if one wishes to infer full information about the channel. We assume knowledge of the state preparations and measurements in the form of matrices in the computational basis. This means the pure entangled state has the form

\[
|\psi\rangle_{AP} = \sum_i s_i |i\rangle_A|p\rangle_P = \sqrt{d_A} \psi_A^{1/2}(\hat{\Phi})_{AP} = \sqrt{d_A} \psi_P^{1/2}(\hat{\Phi})_{AP}.
\]

where \( \beta \) is the identity map. Again, the collection of measurements should be informationally complete if one wishes to infer full information about the channel. We assume knowledge of the state preparations and measurements in the form of matrices in the computational basis. This means the pure entangled state has the form

\[
|\psi\rangle_{AP} = \sum_i s_i |i\rangle_A|p\rangle_P = \sqrt{d_A} \psi_A^{1/2}(\hat{\Phi})_{AP} = \sqrt{d_A} \psi_P^{1/2}(\hat{\Phi})_{AP},
\]

where \( \psi_A, \psi_P \) are the respective reduced states on \( A \) and \( P \) of \( |\psi_{AP}\rangle \langle \psi_{AP}| \) and \( |\hat{\Phi}_{AP}\rangle \langle \hat{\Phi}_{AP}| \) the maximally entangled state on \( \mathcal{H}_{AP} \). Note that not all pure states on \( AP \) have this form, but we assume it without loss of generality by redefining \( |\hat{\Phi}_{AP}\rangle \) if necessary.

The tomography procedure proceeds according to Algorithm I. In each round, we prepare \( |\psi\rangle_{AP} \) and we apply the unknown channel \( \Lambda_{A\rightarrow B} \otimes \mathcal{I}_P \rightarrow P \). We then perform a measurement on the bipartite output system \( BP \) using a setting of our choice, yielding an outcome POVM effect \( E^{(i)}_k \). The data set stores all the outcomes of different rounds. In other words, the ancilla-assisted scheme actually realizes the (theoretical) Choi-Jamiolkowski isomorphism in the laboratory under the assumption of the input state and the channel.

The likelihood function for this scheme is given by

\[
L_{AA}(A|E) = \prod_{k,\ell} \left[ \text{tr} \left( \Lambda_{A\rightarrow B}(\psi_{AP}) E^{(i)}_k \right) \right]^{n_{k,\ell}},
\]

where \( n_{k,\ell} \) is the number of times the POVM effect \( E^{(i)}_k \) appears in the data set \( E \). Since

\[
\Lambda_{A\rightarrow B}(\psi_{AP}) = d_A \psi_A^{1/2} \Lambda_{BP} \psi_P^{1/2},
\]

where \( \Lambda_{BP} \) is the corresponding Choi state, we have

\[
L_{AA}(A|E) = \prod_{k,\ell} \left[ \text{tr} \left( \Lambda_{BP} \psi_P^{1/2} E^{(i)}_k \psi_P^{1/2} \right) \right]^{n_{k,\ell}},
\]

where \( \otimes_{k,\ell} \) ranges over the observed data set \( E \).

Since quantum processes correspond to bipartite quantum states via the Choi-Jamiolkowski isomorphism, we can generalize quantum-state tomography methods to quantum processes. Here, we directly apply the existing procedure of Faist and Renner [22] designed for quantum states to infer information about quantum processes. The main result in this section is Theorem 3. We first recall the procedure of constructing confidence region estimators for quantum states, phrased in terms of bipartite states in anticipation with the connection to quantum processes.

### 1. Christandl-Renner confidence regions

Given access to \( n \) copies of an unknown state \( \rho_{AB} \), we can perform a (joint or collective) POVM measurement on \( \rho_{AB}^{\otimes n} \) and upon receiving the data set \( E \), the Christandl-Renner procedure outputs a distribution

\[
d\mu_E(\sigma_{AB}) := c^{-1} \text{tr} (\sigma_{AB}^{\otimes n} E) d\sigma_{AB},
\]

where
Then, the mapping $E$ distribution $\text{tr}(cE)$ on bipartite density matrices. Confidence regions for the unknown $\rho_{AB}$ can be constructed from $d\mu_E(\sigma_{AB})$ as the following proposition asserts.

**Theorem 1 (of main text).** Let $n$ be the number of systems measured by a POVM during tomography and $1 - \epsilon$ be the desired confidence level. For each effect $E$ in the POVM, let $S_{\mu_E} \subseteq \mathcal{D}(\mathcal{H}_{AB})$ be a set of states such that

$$
\int_{S_{\mu_E}} d\mu_E(\sigma_{AB}) \geq 1 - \epsilon \left(\sum_{n=2}^n s_{n,d}^{2n} d_{n,a}^{-1}\right),
$$

where $s_{n,d} = (n + 1)^{d-1} \leq (n + 1)^d - 1$ and let $S_{\mu_E}$ be the enlargement of $S_{\mu_E}$ defined as

$$
S_{\mu_E}^\delta := \{\sigma_{AB} : \exists \sigma' \in S_{\mu_E} \text{ with } P(\sigma, \sigma') \leq \delta\}.
$$

Then, the mapping $E \mapsto S_{\mu_E}^\delta$ is a confidence region estimator for the unknown $\rho_{AB}$ with confidence level $1 - \epsilon$ if

$$
\delta^2 = \frac{2}{n} \left(\ln \frac{2}{\epsilon} + 2 \ln \sum_{n=2}^n s_{n,d}^{2n} d_{n,a}^{-1}\right).
$$

In other words, for any $\rho_{AB} \in \mathcal{D}(\mathcal{H}_{AB})$,

$$
\Pr_E[\rho_{AB} \in S_{\mu_E}^\delta] \geq 1 - \epsilon,
$$

where the probability is taken over the random data set $E$ with distribution $\text{tr}(\rho_{AB}^E)$.

2. Mapping channel tomography to bipartite-state tomography

Consider the ancilla-assisted scheme. In order to learn what the channel $\Lambda_{A\rightarrow B}$ is, we may carry out the experiment as described in Algorithm I, and use the outcome measurements to perform full tomography on the output state $\rho_{BP} := \Lambda_{A\rightarrow B}(\psi_{AP})$. We may then ask, what does this tell us about the unknown channel $\Lambda_{A\rightarrow B}$?

Observe that if we knew the output state $\rho_{BP}$ exactly (limit of infinite data) and assume the input state $\psi_{AP}$ has full rank, then we could read out the true channel: its Choi state is simply given as $\Lambda_{BP} = d_A^{-1} \rho_{BP}^{1/2} \rho_{BP} \rho_{BP}^{1/2}$. Indeed, we have

$$
\rho_{BP} := \Lambda_{A\rightarrow B}(\psi_{AP}) = d_A^{-1/2} \Lambda_{A\rightarrow B}(\Phi_{AB}) \psi_{AP}^{1/2} = d_A^{-1/2} \Lambda_{BP} \rho_{BP}^{1/2}
$$

since under the assumption that $P$ has undergone identity transformation it follows $\psi_{AP} = \rho_{BP}$. Note that this is the same trick used in Appendix A1c to derive the semidefinite program.

Thanks to this observation, we may use the quantum-state tomography method of Ref. [22] to construct confidence regions on the space of quantum processes $\Lambda_{A\rightarrow B}$, as well as on a figure of merit such as the diamond norm to an ideal channel. To do so, we ignore the knowledge of the exact input state $\psi_{A}$, but we assume the global state $|\psi_{AP}\rangle$ has full Schmidt rank (i.e., forgetting the Schmidt coefficients). Upon observing the data set $E$, the classical data processing returns a bipartite state region $S_{\mu_E}^\delta$, which contains information about the pair $(\Lambda_{A\rightarrow B}, \psi_{A})$. The interpretation of $S_{\mu_E}^\delta$ is given by Theorem 1, and together with the observation above [see Eq. (B10)] we have

$$
\Pr_E[\rho_{BP}^{1/2} \Lambda_{BP} \rho_{BP}^{1/2} \in S_{\mu_E}^\delta] \geq 1 - \epsilon,
$$

where the probability is taken over all possible data set $E$ with distribution $\text{tr}(\rho_{BP}^{1/2} \Lambda_{BP} \rho_{BP}^{1/2})$. To recover information about the channel $\Lambda$, for each $\rho_{BP} \in S_{\mu_E}^\delta$, we apply the (completely positive) transformation $T$ defined as

$$
T : \text{End}(\mathcal{H}_{AB}) \rightarrow \text{End}(\mathcal{H}_{AB}),
$$

$$
\rho_{BP} \mapsto d_A^{-1} \rho_{BP}^{1/2} \Lambda_{BP} \rho_{BP}^{1/2}.
$$

Observe that $T$ maps any $\rho_{BP}$ with full-rank marginal $\rho_{BP}$ to a Choi state. Also, the set $S_{\mu_E}^\delta$ contains only $\rho_{BP}$ with full-rank marginal $\rho_{BP}$ because we only sample according to the uniform measure $d\sigma_{AB}$ (i.e., the set of rank-deficient $\rho_{BP}$ has measure zero). This means the image of $S_{\mu_E}^\delta$ under $T$ will be a set of Choi matrices which can be interpreted via Choi-Jamiolkowski as a region of quantum processes (completely positive and trace-preserving maps). We conclude

$$
\Pr_E[\Lambda_{BP} \in T(S_{\mu_E}^\delta)] \geq 1 - \epsilon,
$$

which implies $T(S_{\mu_E}^\delta)$ are confidence regions for quantum processes.

3. Regions for figures of merit

The confidence region on channel space constructed in the last section contains full information on the unknown channel. But, if one is only interested in a property of the channel, for instance how close is it to an ideal process, then obtaining confidence region for a given figure of merit suffices. We now present how one can do this using push forward of measures.

Given a figure of merit for quantum processes $f_{\text{channel}}$ (defined on channel space), we associate a function $f$ defined on the set of bipartite states as

$$
f(\rho_{BP}) := f_{\text{channel}}(J^{-1}(d_A^{-1/2} \rho_{BP}^{1/2} \rho_{BP}^{1/2})).
$$

which is just $f_{\text{channel}}$ acting on the channel $J^{-1}(d_A^{-1/2} \rho_{BP}^{1/2} \rho_{BP}^{1/2})$ obtained from $\rho_{BP}$ via the mapping $T$. This allows us to directly use the tools of Ref. [22] to obtain confidence intervals for the figure of merit $f$ which will yield the same result as $f_{\text{channel}}$. Explicitly, for any $v \in \mathbb{R}$,

$$
\mu(v) = \int d\mu_E(\sigma_{AB}) \delta(f(\sigma_{AB}) - v)
$$

is the probability density of the push forward of $d\mu_E(\sigma_{AB})$ along $f$. This density provides confidence region for a figure of merit as certified by the following proposition.

**Theorem 3.** Let $\mu_E$ be given as in (B5), and let $\mu(v)$ be defined as in (B15). Then, for any threshold value $\psi_{\text{thres}} > 0$, the region

$$
R_{\text{thres},\delta} = \{\rho_{AB} : f(\rho_{AB}) \leq \psi_{\text{thres}} + O(\delta)\}
$$

is a confidence region for the figure of merit $f$ to $\psi_{\text{thres}}$ with confidence level $1 - \epsilon$. This theorem is in fact, in addition to reconstructing the channel $\Lambda_{A\rightarrow B}$, we may use this procedure to confirm the correct preparation of the input state $\sigma_A$. 

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5In fact, in addition to reconstructing the channel $\Lambda_{A\rightarrow B}$, we may use this procedure to confirm the correct preparation of the input state $\sigma_A$. 

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of states representing channels at least $v_{\text{thres}} + O(\delta)$ close to the reference channel, is a confidence region of confidence level $1 - \epsilon$ where

$$\epsilon = \text{poly}(n) \left[ 1 - \int_0^{v_{\text{thres}}} \mu(v) dv \right]. \quad (B17)$$

In summary, for ancilla-assisted tomography scheme, determining the histogram $\mu(v)$ in (B15) gives us all the necessary information to construct confidence regions of any confidence level in terms of the figure of merit $f(\text{channel}(\Lambda_{A\rightarrow B})$.

$$f(\rho_{BP}) = \frac{1}{2} \max \left\{ \left| \rho_p^{-1/2} \rho_{BP} \rho_p^{-1/2} - d_A^\text{ideal}_{BP} \right|_1 \right\}, \quad (B18)$$

One is left to perform a numerical computation of $\mu(v)$ for the above function $f$, as explained in details in Ref. [22].

**APPENDIX C: CHANNEL-SPACE SAMPLING METHOD**

This method applies to either the ancilla-assisted scheme explained in the previous Appendix, or the prepare-and-measure scheme where no entanglement is required. In the prepare-measure scheme [see Fig. 2(a)], we select a collection of input states $\sigma^i$, and select a collection of measurements $E^{(i)} = \{E_{k}^{(i)}\}$. This set of state preparation and measurement (SPAM) should be informationally complete if one wishes to fully reconstruct the unknown channel. The SPAM is represented as certain set of matrices in the computational basis $|i\rangle : i = 0, \ldots, d_A - 1$.

The data collection procedure goes as follows: in each round, we choose an input state $\sigma^i$, we choose a measurement $\ell$ on output, we send $\sigma^j$ through the channel, and record the measurement outcome $k$ on the output. The data set $E$ consists of all pairs $(\sigma^i, E_{k}^{(i)})$ chosen and observed for each round.

Typically, one can choose the states $j$ in order, i.e., first perform measurements on $\sigma^1$, then on $\sigma^2$, etc. The choice of the output measurement setting $\ell$ is allowed to depend on $j$. Since we are under i.i.d. channel assumption, at each round it is the same unknown channel $\Lambda$ which is applied, and that previous outcomes have no influence on new rounds.

The likelihood function for a data set $E$ in this scenario is defined using the matrix representations of the SPAM according to Born’s rule

$$L_{PM}(\Lambda|E) = \prod_{j,k,\ell} \left[ \text{tr}(A_{\Lambda}(\sigma^j) E_{k}^{(i)}) \right]^{n_{j,k,\ell}}, \quad (C1)$$

where $n_{j,k,\ell}$ is the number of times the given pair $(\sigma^j, E_{k}^{(i)})$ appears in the data set $E$. Using (A3), we rewrite the likelihood function as

$$L_{PM}(\Lambda|E) = \prod_{j,k,\ell} \left[ d_A \text{tr}(A_{\Lambda AB}(\sigma^j)^\dagger \otimes E_{k}^{(i)}) \right]^{n_{j,k,\ell}} = d_A^n \left( A_{\Lambda AB}^{\otimes n} \otimes (\sigma^j)^\dagger \otimes E_{k}^{(i)} \right), \quad (C2)$$

where $\otimes_{j,k,\ell}$ ranges over the observed data set $E$.

**Diamond distance to ideal and worst-case entanglement fidelity**

The methodology outlined in the previous paragraphs can be specialized to the diamond distance to an ideal reference channel $\Lambda_{\text{ideal}}$. Here, we take

$$f(\text{channel}(\Lambda_{B\rightarrow p})) = f_p(\Lambda_{B\rightarrow p}) = \frac{1}{2} \left\| \Lambda_{B\rightarrow p} - \Lambda_{\text{ideal}}_{B\rightarrow p} \right\|_1 \quad (B18)$$

to be the desired figure of merit on channel space. This induces a figure of merit in the space of bipartite quantum states

$$\rho, W : W \leq 1_B \otimes \bar{\rho}, W \geq 0, \bar{\rho} \in D(\mathcal{X}), \quad (B19)$$

The method in the previous section maps a channel tomography problem into a (constrained) bipartite-state tomography problem. One may ask if this is the only solution. In this section, we provide an alternative construction natively on the channel space. This has consequence on the numerical implementation: we no longer need to sample from bipartite-state space. Instead, we can directly sample “random channels” which leads to improved numerical efficiency. The main results in this section are Theorems 2 and 4.

1. Regions on channel space

Inspired by the Christandl-Renner construction [21], we define the following confidence region estimator for quantum processes. Our confidence region is constructed from the probability measure on the space of quantum processes $\mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B)$:

$$d_{\mathcal{V}_\Lambda}(\Lambda) := c_{\mathcal{V}_\Lambda}^{-1} \mathcal{L}(\Lambda|E) d\nu(\Lambda), \quad (C3)$$

where $\mathcal{L}(\Lambda|E)$ is either prepare-measure or ancilla-assisted likelihood and $c_{\mathcal{V}_\Lambda} = \int \mathcal{L}(\Lambda|E) d\nu(\Lambda)$ serves as a normalizing constant and $d\nu(\Lambda)$ is the induced measure on $\mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B)$ defined in Appendix A 2.

The main result in this section is as follows:  
*Theorem (2 of main text).* Let $n$ be the number of channel uses during tomography and $1 - \epsilon$ be the desired confidence level. For each data set $E$, let $R_{\epsilon} \subseteq \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B)$ be a set of channels such that

$$\int_{R_{\epsilon}} d_{\mathcal{V}_\Lambda}(\Lambda) \geq 1 - \epsilon \frac{s_{n,d}^2}{2n^d d_{\epsilon}}, \quad (C4)$$

where $s_{n,d} \leq \binom{n+d-1}{d-1} \leq (n+1)^{d-1}$ and let $R_{\epsilon}$ be the enlargement

$$R_{\epsilon} := \{ \Lambda \in \mathcal{C}(\mathcal{H}_A \rightarrow \mathcal{H}_B) : \forall \Lambda' \in R_{\epsilon} \} \subseteq H \Lambda, \Lambda' \in \delta \}, \quad (C5)$$

Then, the mapping $E \mapsto R_{\epsilon}$ is a confidence region estimator for the unknown $\Lambda_{A \rightarrow B}$ with confidence level $1 - \epsilon$ if

$$\delta^2 = \frac{2}{n} \ln \frac{2}{\epsilon} + 3 \ln s_{2n,d_{\epsilon}}. \quad (C6)$$
In other words, for all channel $\Lambda \in \mathcal{C} (\mathcal{H}_A \rightarrow \mathcal{H}_B)$

$$\Pr_E [\Lambda_{A \rightarrow B} \in \mathcal{R}_{\nu}^E] \geq 1 - \epsilon,$$  \hfill (C7)

where the probability is over the random data set $E$ with distribution $\Pr (E | \Lambda) = \mathcal{L} (\Lambda | E)$.

Before starting the proof, we will need the following results.

Proposition 1. For any channel $\Lambda_{A \rightarrow B}$, if $| \Lambda \rangle \in \mathcal{H}_{A B A B'}$ is a purification of its Choi state, then

$$| \Lambda \rangle | \Lambda \rangle^{\otimes n} \leq s_n^2 \zeta \int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle^{\otimes n}$$

$$= s_n^2 \zeta \int dU U_{BA B'}^\otimes | \Psi_0 \rangle^{\otimes n} U_{BA B'}^{\dagger \otimes n},$$ \hfill (C8)

where $s_n := (n + d - 1)$. \hfill Proof

The main idea of this proof is to discretize the Haar integral using Caratheodory’s theorem, and dominate the left-hand side by a trivial operator inequality. By definition, the operator

$$\int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle^{\otimes n}$$

lies in the convex hull of the set $\{ | \Psi \rangle | \Psi \rangle^{\otimes n} : | \Psi \rangle \in \mathcal{S} \mathcal{E} \}$, whose linear span (in the ambient space $\text{End} (\mathcal{H}_{A A B B'})$) has dimension $D$. By Caratheodory’s theorem, there exists a convex combination $(q_i, | \Psi_i \rangle | \Psi_i \rangle^{\otimes n})$ with size at most $D + 1$ such that

$$\int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle^{\otimes n} = \sum_{i=1}^{D+1} q_i | \Psi_i \rangle | \Psi_i \rangle^{\otimes n},$$ \hfill (C9)

Among the probability weights $q_i$ there exists a largest element denoted $q_{\max}$ and its associated purified Choi state $| \Psi_{\max} \rangle | \Psi_{\max} \rangle^\otimes n$, from which we split off this term in the finite sum as

$$\sum_{i=1}^{D+1} q_i | \Psi_i \rangle | \Psi_i \rangle^{\otimes n} = q_{\max} | \Psi_{\max} \rangle | \Psi_{\max} \rangle^{\otimes n} + \sum_{i=\max} \sum_{i=\max} q_i | \Psi_i \rangle | \Psi_i \rangle^{\otimes n}.$$ \hfill (C10)

By left invariance of the measure $d\nu (| \Psi \rangle)$ and the (unitary) structure of the set $\mathcal{S} \mathcal{E}$, we can without loss of generality assume that $\Psi_{\max} = \Lambda$. More precisely, let $W_{A B B'}$ be a unitary transformation bringing $| \Psi_{\max} \rangle$ to $| \Lambda \rangle$, we have (leaving the system label $B A B'$ implicit)

$$W^{\otimes n} \left( \int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle | \Psi \rangle^{\otimes n} \right) W^{\otimes n}$$

$$= \sum_{i=\max} q_i W^{\otimes n} | \Psi_i \rangle | \Psi_i \rangle^{\otimes n} W^{\otimes n}.$$ \hfill (C12)

Using linearity of integration and translational invariance of the integrating measure, this equation simplifies to

$$\int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle | \Psi \rangle^{\otimes n}$$

$$= q_{\max} | \Lambda \rangle | \Lambda \rangle^{\otimes n} + \sum_{i=\max} q_i | \Psi_i \rangle | \Psi_i \rangle^{\otimes n},$$ \hfill (C13)

where $| \Psi_i \rangle$ is some other vector in $\mathcal{S} \mathcal{E}$.

Now, since all operators in the convex combination are positive semidefinite, we obtain

$$\int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle | \Psi \rangle^{\otimes n} \geq q_{\max} | \Lambda \rangle | \Lambda \rangle^{\otimes n}.$$ \hfill (C14)

By the property of the maximum weight $q_{\max}$, namely, $q_{\max} \geq 1 / (D + 1)$, we get

$$| \Lambda \rangle | \Lambda \rangle^{\otimes n} \leq (D + 1) \int_{\mathcal{S} \mathcal{E}} d\nu (| \Psi \rangle) | \Psi \rangle^{\otimes n}.$$ \hfill (C15)

Finally, span$\{ | \Psi \rangle | \Psi \rangle^{\otimes n} : | \Psi \rangle \in \mathcal{S} \mathcal{E} \} \subseteq \text{span} \{ | \Psi \rangle | \Psi \rangle^{\otimes n} : | \Psi \rangle \in \mathcal{H}_{A B A B'} \}$ and the latter is identified as a subspace of $\text{End} (\mathcal{H}_{A B A B'})$, the operator space on the symmetric subspace of $\mathcal{H}_{A B A B'}^\otimes n$. Together with the constraint that trace is 1, we have $D \leq s_{n,d}(\mathcal{A} B) - 1$ where the dimension of the symmetric subspace is $s_{n,d} := (n + d - 1)$. This completes the proof of the operator inequality. \hfill \blacksquare

Proof of Theorem 2. Our proof technique follows closely that of [21], with the main technical difficulty being incorporating the a priori constraint $\text{tr}_{B} (\Lambda_{A B}) = 1$. This allows the reduction of numerical sampling from bipartite state to channel space. For any region estimator, our construction $E \mapsto R_{\nu}^E$, in particular, the failure probability of the reconstruction typically depends on the underlying unknown channel

$$P_{\text{fail}} (\Lambda_{A \rightarrow B}) = \Pr_E [\Lambda_{A \rightarrow B} \notin \mathcal{R}_{\nu}^E] = \sum_{E} \Pr (E | \Lambda) \chi (\Lambda_{A \rightarrow B} ; \mathcal{R}_{\nu}^E),$$ \hfill (C16)

where $\Pr (E | \Lambda)$ is the probability of obtaining data set $E$, and $\chi (\Lambda_{A \rightarrow B} ; \mathcal{R}_{\nu}^E)$ is the indicator function of the set $\mathcal{R}_{\nu}^E := \mathcal{C} (\mathcal{H}_A \rightarrow \mathcal{H}_B) \setminus \mathcal{R}_{\nu}^E$ (i.e., the complement set). Recall that

$$\Pr (E | \Lambda) = \begin{cases} d_n^2 \text{tr} (\Lambda_{AB}^{\otimes n} \otimes_{\nu,k} (\sigma_{\chi}^T \otimes E^{(\ell)}), & \text{in prepare-and-measure scheme}, \\ d_n^2 \text{tr} (\Lambda_{AB}^{\otimes n} \otimes_{\nu,k} \psi_{\nu}^{(\ell)}/2 \otimes E^{(\ell)}), & \text{in ancilla-assisted scheme}. \end{cases}$$ \hfill (C17)

Our goal will be bounding this failure probability independently of $\Lambda_{A \rightarrow B}$ by using the operator inequality we have just developed.
POVM with effects $T$ and $1 - T$ acting on the quantum state $|\Lambda\rangle_{ABA'B'}$. Upon the observation of the effect $T$, we decide that $|\Lambda\rangle_{ABA'B'} \in Q_{v_e}^T$ and similarly for $1 - T$. In other words, we are approximating $\chi(|\Lambda\rangle_{ABA'B'}; Q_{v_e}^e)$ by a quantum measurement. Here, we construct such an approximation using Holevo’s covariant measurement \[49\].

Let $k$ be the number of copies of $|\Lambda\rangle \in H_{ABA'B'}$ used in the approximation, i.e., we are given $|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle$. If we ignore the fact that $|\Lambda\rangle \in \mathcal{P}^e$, we can use the Holevo’s continuous POVM $\{s, d, \nu\}$ to distinguish $|\Lambda\rangle \in H_{ABA'B'}$ among the set of pure states. Here, $d\nu$ is the uniform spherical measure on the set of pure states of $H_{ABA'B'}$ and $s, d, \nu$ is the dimension of the symmetric subspace of $\mathcal{H}_{ABA'B'}$. Coarse graining this measurement, we can distinguish $|\Lambda\rangle_{ABA'B'} \in Q_{v_e}^T$ versus $|\Lambda\rangle_{ABA'B'} \in Q_{v_e}^e$ by the following POVM with two effects (analogous to Ref. [21]):

\[
T_{v_e}^{1/2} := s_{k,d^2} \int_{Q_{v_e}^T} |\phi\rangle \langle \phi|^R d\phi, \quad 1 - T_{v_e}^{1/2}.
\] (C21)

We now check that this POVM indeed approximates $\chi(|\Lambda\rangle; Q_{v_e}^e)$. For all $|\Lambda\rangle \in \overline{Q}_{v_e}^e$, using the definition of $\chi(|\Lambda\rangle; Q_{v_e}^T)$,

\[
\chi(|\Lambda\rangle; Q_{v_e}^T) - \text{tr}(|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle T_{Q_{v_e}^T}) = 1 - s_{k,d^2} \int_{Q_{v_e}^T} \text{tr}(|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle \langle \phi| \langle \phi|^R d\phi.
\] (C22)

Since $|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle$ is supported on the symmetric subspace, we reinterpret the constant 1 above as

\[
1 = \text{tr}(|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle s_{k,d^2} \int |\phi\rangle \langle \phi|^R d\phi.
\] (C23)

which implies for all $|\Lambda\rangle \in \overline{Q}_{v_e}^e$

\[
\chi(|\Lambda\rangle; Q_{v_e}^T) - \text{tr}(|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle T_{Q_{v_e}^T}) = s_{k,d^2} \int_{Q_{v_e}^T} \text{tr}(|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle \langle \phi| \langle \phi|^R d\phi.
\] (C24)

By the definition of the sets

\[
R_{v_e}^{1/2} := \{ \Psi \in \mathcal{C}(H_{ABA'}) : \exists \Psi' \in R_{v_e} \text{ with } P(\Psi, \Psi') \leq \delta/2 \}
\] (C27)

and

\[
\overline{R}_{v_e}^{1/2} := \mathcal{C}(H_{ABA'}) \setminus \{ \Psi \in \mathcal{C}(H_{ABA'}) : \exists \Psi' \in R_{v_e} \text{ with } P(\Psi, \Psi') \leq \delta \},
\] (C28)

we have, for $\Lambda_{AB} \in \overline{R}_{v_e}^{1/2}$ and $\phi_{AB} := \text{tr}_{AB}|\phi\rangle \langle \phi| \in R_{v_e}^{1/2},$

\[
F(\Lambda_{AB}, \phi_{AB}) = \sqrt{1 - P(\Lambda_{AB}, \phi_{AB})^2} \leq \sqrt{1 - (\delta/2)^2} \leq e^{-\delta^2/2},
\] (C29)

using the reverse triangle inequality for purified distance. In summary, we obtain the approximation

\[
\chi(|\Lambda\rangle; Q_{v_e}^T) - \text{tr}(|\Lambda\rangle|\Lambda\rangle|\Lambda\rangle T_{Q_{v_e}^T}) \leq \epsilon_1 := s_{k,d^2} e^{-\delta^2/2}.
\] (C30)
Now, we can start bounding the failure probability. Inserting (C30) into (C19), we have an intermediate bound
\[
P_{\text{fail}}(\Lambda_{AB}) \leq \epsilon_1 + \sum_E d_E^* \text{tr}\left[|\Lambda\rangle\langle\Lambda|_{\Lambda_{AB}} \rho_{\Lambda_{AB}} \otimes E_{\Lambda} \right] \text{tr}\left[|\Lambda\rangle\langle\Lambda|^\otimes_k T_{Q_{\epsilon/2}}\right]
\]
\[
= \epsilon_1 + \sum_E d_E^* \text{tr}\left[|\Lambda\rangle\langle\Lambda|^\otimes_k \rho_{\Lambda_{AB}} \otimes E_{\Lambda} \otimes T_{Q_{\epsilon/2}}\right].
\]
(C31)

Using the operator inequality in the Proposition 1, namely,
\[
|\Lambda\rangle\langle\Lambda|^\otimes_k \leq s_{n+k,d_{\Lambda}}^2 \int_{\mathcal{P}(\mathcal{E})} d\nu(|\Psi\rangle) |\Psi\rangle\langle\Psi|^\otimes_k,
\]
we can bound the right-hand side independent of the unknown \(\Lambda_{AB}\) as follows:
\[
P_{\text{fail}}(\Lambda_{AB}) \leq \epsilon_1 + s_{n+k,d_{\Lambda}}^2 \int_E d\nu(|\Psi\rangle) d_E^* \text{tr}\left[|\Psi\rangle\langle\Psi|^\otimes_k \rho_{\Lambda_{AB}} \otimes E_{\Lambda} \otimes T_{Q_{\epsilon/2}}\right]
\]
\[
= \epsilon_1 + s_{n+k,d_{\Lambda}}^2 \sum_E c_E^* \int_E d\nu_E(|\Psi\rangle)\text{tr}\left[|\Psi\rangle\langle\Psi|^\otimes_k T_{Q_{\epsilon/2}}\right].
\]
(C32)

where the last equality follows from the definition of the a posteriori measure \(d\nu_E(\Psi)\). For each measurement outcome \(E\), the integral can split into two parts based on the set \(R_{\nu_E}\) from which the kernels are uniformly bounded as follows:
\[
\int_{R_{\nu_E}} d\nu_E(|\Psi\rangle)\text{tr}\left[|\Psi\rangle\langle\Psi|^\otimes_k T_{Q_{\epsilon/2}}\right] \leq s_{k,d_{\nu_E}} 2^{k/2} \leq s_{k,d_{\nu_E}} e^{-k^2/2},
\]
(C33)

using the definition of \(T_{Q_{\epsilon/2}}\) and the fidelity bound \(F(\Psi_{AB} \in R_{\nu_E}, \Phi_{AB} \in R_{\nu_E}^k) \leq \sqrt{1 - (\delta/2)^2}\), and
\[
\int_{R_{\nu_E}} d\nu_E(|\Psi\rangle)\text{tr}\left[|\Psi\rangle\langle\Psi|^\otimes_k T_{Q_{\epsilon/2}}\right] \leq \int_{R_{\nu_E}} d\nu_E(|\Psi\rangle)
\]
(C34)

since \(\text{tr}(|\Psi\rangle\langle\Psi|^\otimes_k T_{Q_{\epsilon/2}}) \leq 1\). Choose \(k = n\), the fact that \(\sum_E c_E^* \leq 1\), and combining all the inequalities together we have
\[
P_{\text{fail}}(\Lambda_{AB}) \leq \epsilon_1 + s_{n+k,d_{\Lambda}}^2 \sum_E c_E^* \int_{R_{\nu_E}} d\nu_E(|\Psi\rangle)
\]
\[
= s_{n,d_{\Lambda}} e^{-n^2/2} + s_{2n,d_{\Lambda}}^2 s_{n,d_{\Lambda}} e^{-n^2/2} + s_{2n,d_{\Lambda}}^2 \sum_E c_E^* \int_{R_{\nu_E}} d\nu_E(|\Psi\rangle)
\]
\[
\leq s_{2n,d_{\Lambda}}^2 e^{-n^2/2} + s_{2n,d_{\Lambda}}^2 \sum_E c_E^* \int_{R_{\nu_E}} d\nu_E(|\Psi\rangle).
\]
(C35)

If we choose \(R_{\nu_E}\) and \(\delta\) such that
\[
\int_{R_{\nu_E}} d\nu_E(|\Psi\rangle) \geq 1 - \epsilon/2 + \epsilon/2 = \epsilon
\]
(C36)

then \(P_{\text{fail}}(\Lambda_{AB}) \leq \epsilon/2 + \epsilon/2 = \epsilon\) as desired. The proof of the Proposition is complete.

2. Regions for figures of merit

The construction of confidence region on channel space can be pushed forward to obtain confidence regions for any figure of merit of channels we are interested in. The idea is exactly the same as [22] and we include it here for completeness. Let \(f_{\text{channel}} : \mathcal{H}_A \rightarrow \mathcal{H}_B \rightarrow \mathbb{R}\) be an arbitrary figure of merit of channels. The measure \(d\nu_E(\Lambda)\) can be pushed forward by \(f_{\text{channel}}\) to a measure on \(\mathbb{R}\), which can then be represented as a density function \(h(v)\) with respect to the Lebesgue measure of \(\mathbb{R}\). Concretely, we have
\[
h(v) = \int d\nu_E(\Lambda) \delta(f_{\text{channel}}(\Lambda) - v),
\]
(C42)

where \(\delta(f_{\text{channel}}(\Lambda) - v)\) is the Dirac delta measure on \(\mathbb{R}\) at the point mass \(v \in \mathbb{R}\). And for some subset of values \(V\), the measure of \(V\) is given by
\[
\int_{f_{\text{channel}}^{-1}(V)} d\nu_E(\Lambda) = \int_V h(v) dv,
\]
(C43)

where \(dv\) is the Lebesgue measure on \(\mathbb{R}\). The density \(h(v)\) allows us to construct confidence interval for the property we desired.
**Proposition 2.** Let \( f_{\text{channel}} \) be a figure of merit and choose a confidence level \( 1 - \epsilon \). For each data set \( E \), let \( V_{\epsilon E} \subseteq \mathbb{R} \) be a region of values such that

\[
\int_{V_{\epsilon E}} h(v)dv \geq 1 - \epsilon, \quad (C44)
\]

and let \( V_{\epsilon E}^\perp \) be defined as

\[
V_{\epsilon E}^\perp := \{ v \in \mathbb{R} : \exists v' \in V_{\epsilon E} \text{ with } |v - v'| \leq \omega_f(\delta) \}, \quad (C45)
\]

where \( \omega_f(\delta) := \sup_{\alpha \in \mathcal{R}_{\Lambda_{\alpha}}} |f(\Lambda) - f(\Lambda')| \). Then, the mapping \( E \mapsto V_{\epsilon E}^\perp \) is a confidence region estimator for the figure of merit \( f_{\text{channel}} \) with confidence level \( 1 - \epsilon \) if

\[
\delta^2 = \frac{2}{n} \left( \ln \frac{2}{\epsilon} + 3 \ln s_{2n,d_{\text{in}}^2} \right). \quad (C46)
\]

In other words, for all channel \( \Lambda \in \mathcal{C}(\mathcal{H}_A \to \mathcal{H}_B) \),

\[
\Pr[ f_{\text{channel}}(\Lambda) \in V_{\epsilon E}^\perp ] \geq 1 - \epsilon. \quad (C47)
\]

**Proof.** It is clear from the fact that, as defined, \( V_{\epsilon E}^\perp \supseteq f_{\text{channel}}(\mathcal{C}_{\perp}(V_{\epsilon E}^\perp)) \).

For each figure of merit of interest, we can derive a bound on \( \omega_f(\delta) \) by simple inequalities for distance measures. For diamond distance, we have the following result.

**Proposition 3.** For each data set \( E \), let \( \gamma_E \in [0, 1] \) be such that

\[
\int_0^{\gamma_E} h(v)dv \geq 1 - \epsilon, \quad (C48)
\]

Then, the mapping \( E \mapsto [0, \gamma_E + d_\Lambda \delta/2] \) is a confidence region estimator for the diamond distance to ideal with confidence level \( 1 - \epsilon \) if

\[
\delta^2 = \frac{2}{n} \left( \ln \frac{2}{\epsilon} + 3 \ln s_{2n,d_{\text{in}}^2} \right). \quad (C49)
\]

In other words, for all channel \( \Lambda \in \mathcal{C}(\mathcal{H}_A \to \mathcal{H}_B) \),

\[
\Pr[ f_{\text{ideal}}(\Lambda_{AB} - \Lambda_{AB}') \leq \gamma_E + d_\Lambda \delta/2 ] \geq 1 - \epsilon, \quad (C50)
\]

\[
|f(\Lambda) - f(\Lambda')| = |f_{\text{worst}}(\Lambda) - f_{\text{worst}}(\Lambda')| \leq \|\Phi(\rho_A(d_\Lambda \Lambda_{AB}))\rho_A(\Phi) - \langle \Phi|\rho_A(d_\Lambda \Lambda_{AB})\rho_A|\Phi\rangle| \leq d_\Lambda \gamma_E \|\Phi(\rho_A(\Lambda_{AB} - \Lambda_{AB}')\rho_A)\Phi\rangle\langle \Phi|\rho_A(\Lambda_{AB} - \Lambda_{AB}')\rho_A|\Phi\rangle \leq d_\Lambda \gamma_E \|\Phi(\rho_A(\Lambda_{AB} - \Lambda_{AB}')\rho_A)\Phi\rangle\langle \Phi|\rho_A(\Lambda_{AB} - \Lambda_{AB}')\rho_A|\Phi\rangle \leq d_\Lambda, \quad (C57)
\]

using Holder inequality for Schatten norms and \( |\Lambda_{AB} - \Lambda_{AB}'| \leq P(\Lambda_{AB}, \Lambda_{AB}') \).

**APPENDIX D: METROPOLIS-HASTINGS ALGORITHM IN CHANNEL SPACE**

The previous two sections describe the construction of confidence region estimators for quantum processes, which utilize distributions \( d\mu_E(\sigma) \) and \( d\nu_E(\Lambda) \). We now describe how one can numerically estimate such distributions so that the densities \( \mu(v) \) and \( h(v) \) can be approximated. The distribution \( d\mu_E(\sigma) \) or the density \( \mu(v) \) can be estimated by numerically producing a lot of samples. These can be generated by the Metropolis-Hastings random walk in (bipartite) state space, where the probability is over the random data set \( E \) with distribution \( \Pr(E|\Lambda) = \mathcal{L}(\Lambda|E) \).

**Proof.** Continuing from the previous Proposition, we set \( V_\epsilon := [0, \gamma_E] \); it remains for us to obtain a bound on \( \omega_f(\delta) \). Using the reverse triangle inequality and semidefinite programming (SDP) reformulation of diamond norm, we have

\[
|f_\epsilon(\Lambda) - f_\epsilon(\Lambda')| = \frac{1}{2} \| \Lambda - \Lambda_{\text{ideal}} \|_1 - \| \Lambda' - \Lambda_{\text{ideal}} \|_1 \leq \frac{1}{2} \| \Lambda_{AB} - \Lambda_{AB}' \|_1, \quad (C51)
\]

where the last inequality utilizes the duality between Schatten 1-norm and Schatten \( \infty \)-norm to bound the objective function of the diamond norm SDP. Since the purified distance dominates the trace distance, we obtain

\[
\frac{1}{2} \| \Lambda_{AB} - \Lambda_{AB}' \|_1 \leq \frac{1}{2} \Pr(\Lambda_{AB} = \Lambda_{AB}'), \quad (C53)
\]

which implies \( \omega_f(\delta) \leq d_\Lambda \delta/2 \).

For worst-case entanglement fidelity, we have the following result.

**Proposition 4.** For each data set \( E \), let \( \gamma_E \in [0, 1] \) be such that

\[
\int_0^{\gamma_E} h(v)dv \geq 1 - \epsilon, \quad (C48)
\]

Then, the mapping \( E \mapsto [0, \gamma_E - d_\Lambda \delta] \) is a confidence region estimator for the diamond distance to ideal with confidence level \( 1 - \epsilon \) if

\[
\delta^2 = \frac{2}{n} \left( \ln \frac{2}{\epsilon} + 3 \ln s_{2n,d_{\text{in}}^2} \right). \quad (C55)
\]

In other words, for all channel \( \Lambda \in \mathcal{C}(\mathcal{H}_A \to \mathcal{H}_B) \),

\[
\Pr[ f_{\text{worst}}(\Lambda_{AB}) \geq \gamma_E - d_\Lambda \delta ] \geq 1 - \epsilon, \quad (C56)
\]

where the probability is over the random data set \( E \) with distribution \( \Pr(E|\Lambda) = \mathcal{L}(\Lambda|E) \).

**Proof.** We set \( V_\epsilon := [\gamma_E, 1] \). Let \( \rho_A \) be an optimizer of \( f_{\text{worst}}(\Lambda') \) since \( \rho_A \) will give an upper bound on \( f_{\text{worst}}(\Lambda) \) we have \( f = f_{\text{worst}} \). whose details can be found in Ref. [22]. Here, we only discuss the Metropolis-Hastings random walk in channel space. Recall that in the channel-space method, we need to be able to compute the density \( h(v) \) for the given figure of merit \( f_{\text{channel}} \). We do this numerically using Metropolis-Hastings algorithm. The output of this algorithm is a histogram of the figure of merit which approximates the continuous density. Let us recall the Metropolis-Hastings algorithm for continuous sample space [38]. Let \( p(x)dx \) be the target distribution from which we want to sample, and \( q(x'|x)dx' \) be a proposal distribution, all displayed with respect to the same base measure \( dx = dx' \). We assume that the proposal density function is symmetric \( q(x'|x) = q(x|x') \). When the process is at point \( x \), the distribution \( q(x'|x)dx' \) proposes a new point \( x' \). If
We want to generate samples from the target distribution
\[
dv_E(\Lambda) := c_E^{-1} L(\Lambda|E) dv(\Lambda),
\]
where \( L(\Lambda|E) \) is the prepare-and-measure or ancilla-assisted likelihood function and \( dv(\Lambda) \) is the induced measure on channel space. Recalling the definition of \( dv(\Lambda) \), we thus want to sample from
\[
dv_{UBA}(U_{BA'B}) := c_{UBA}^{-1} L(U_{BA'B}|E) dv_{UBA},
\]
with \( dv_{UBA} \) the invariant Haar measure. Concretely, in the prepare-and-measure scheme we take
\[
\mathcal{L}_{PM}(U|E) = d^{n}_{BA'} \left( (U|\Psi_0\rangle\langle\Psi_0|U^{-1}) \bigotimes_{j,k,\ell} (\sigma_j^i)^{T} \otimes E_k^{(i)} \right)
\]
and in the ancilla-assisted scheme we take
\[
\mathcal{L}_{AA}(U|E) = d^{n}_{tr} \left( (U|\Psi_0\rangle\langle\Psi_0|U^{-1}) \bigotimes_{k,\ell} \psi_k^{1/2} E_k^{(i)} \psi_k^{1/2} \right).
\]
where \( |\Psi_0\rangle \) is the fixed reference state in (A15). This can be done using the Metropolis-Hastings algorithm, by designing a symmetric proposal distribution over the space of all unitaries \( U_{BA'B} \), and setting \( q(U_{BA'B}|U_{BA'}) \propto L(U_{BA'}|E) \). To ensure \( q(U'|U) = q(U|U') \), let \( q(W)dw \) be a distribution on unitaries on \( BA'B \) such that \( q(W) = q(W^{-1}) \). For each point \( U \), if we define \( U' := WU \), then we have a symmetric proposal distribution \( q(U'|U) = q(WU|U) = q(W) = q(W^{-1}) = q(W^{-1}U'|U') = q(U|U') \), namely, \( q(WU|U)dw \) where \( dw \) is the Haar measure. It remains to fix \( q(W)dw \) with \( q(W) = q(W^{-1}) \). We have implemented two choices:

(i) “\( e^{iH} \)-type jumps”: We pick a random \( d_{BA'B} \times d_{BA'B} \) matrix \( N \) with each entry independent and normally distributed complex numbers with standard deviation given by the step size. We then calculate \( H = N + N^\dagger \) and set \( W = e^{iH} \), inducing a measure \( q(W)dw \). Denoting by \( dN \) the measure induced on \( N \) by this sampling procedure, observe that \( dN = d(-N) \) as the normal distribution is symmetric.

Furthermore, the Haar measure is invariant under the adjoint, \( dw = d(W) \), since \( d(W) \) is also unitarily invariant and is thus also the Haar measure. Hence, \( q(W)dw = dN = d(-N) = q(W^{-1})d(W') = q(W^{-1})dw \) as required.

(ii) “Elementary rotation jumps”: Choose \( m \in \{x, y, z\} \) uniformly at random and choose two indices \( i < j \) uniformly at random. Choose \( \sin(\alpha) \) at random (normally distributed number whose standard deviation is the step size; truncated to \([-1, 1]) \). Define the unitary \( W_i \) as the qubit rotation on the subspace spanned by \([i, j]|j\rangle \rangle \) defined by \( e^{i\alpha \sigma_m \cdot \sigma_j} \). We see that \(-\alpha (e_m \cdot \sigma)\) is sampled with the same probability as \( \alpha (e_m \cdot \sigma) \) and hence for the same reason as above, \( q(W) = q(W^\dagger) \). In order to keep the acceptance ratio at a reasonable rate, we sample \( N_{\text{num-iter}} \) different instances of \( W_i \), and multiply them together to form the sampled \( W \). One should choose \( N_{\text{num-iter}} \) such that it is possible to keep the acceptance rate around 30%.

### APPENDIX E: IMPACT OF THE RANDOM WALKER AND EMPIRICAL FIT FUNCTION

Overall, we have two main methods, bipartite-state and channel-space sampling, for tomography of an unknown quantum process. In each of these methods, there are several options on how to configure the Metropolis-Hastings random walk: standard or optimized jump distribution for bipartite-state sampling, and eiH or elementary rotation jump distribution for channel-space sampling. The histogram bins are further subjected to a statistical fit analysis to the default fit function. Thus, in this section we study the variability of our simulation results, briefly presented in the main text, in terms of these choices. The main takeaway message is that difference choices lead to very consistent results although their computational performance (time, memory) may differ significantly.

Let us focus on the two-qubit quantum memory example where a simulated data set of \( n = 40,500 \) measurement outcomes is subjected to different numerical analyses. The running times vary quite significantly. For example, the bipartite-state method runs for 85 min with standard jump distribution and 18 min with optimized jump distribution to compute diamond norm distance. The channel-space method runs for 15 h with eiH jump distributions and 2 h with elementary rotation jump distribution. Despite the difference in running times, the histograms produced are consistent across different choices (see Fig. 6, left plot).

The histograms can be compactly described by the fit parameters (or equivalently the quantum error bars) associated with the default fit function. As seen in the main text, the default fit function could yield high goodness of fit or small reduced \( \chi^2 \) values. In the right plot of Fig. 6, we show a situation where the default fit function does not fit well to simulated data. This leads us to develop an empirical model

\[
\ln \mu_{\text{fit}}^{\#2}(v) = -a_2 v^2 - a_1 v + m (\ln v)^p + c.
\]

Trying the fit \#2 yields better agreement and smaller \( \chi^2 \) values. The main lesson here is that the fit models need to be treated a bit delicately on a case-by-case basis. Since good fit is usually achieved for a region of high probability density, we expect that this translates to reasonably accurate confidence region computations. However, in a fully paranoid setting, one might need to be more conservative and use an optimized fit function before computing confidence regions.

### APPENDIX F: CONVERGENCE IN NUMBER OF SAMPLES \( N \rightarrow \infty \)

We now turn to an example where we clearly observe the convergence of the distributions \( h(v) \) and \( \mu(v) \) around the known true figure of merit. Consider a noisy identity process
on a qutrit, of the form
\[ \Lambda_{A\rightarrow B}(\rho) = p\, \rho + (1-p)\, d_B^{-1}\, \mathbb{I}_B, \]  
with \( p = 0.96 \) and \( d_B = 3 \). This gives us a diamond norm to the identity process of
\[ \frac{1}{2} |\Lambda_{A\rightarrow B} - \mathbb{I}_{A\rightarrow B}|_\diamond = 0.035 56. \]

We consider measurements on the input and output systems given by using the Gell-Mann matrices as observables:

\[
\begin{align*}
\lambda_1 &= \begin{pmatrix} 0 & 1 & 0 \\
1 & 0 & 0 \\
0 & 0 & 0 \end{pmatrix}; & \lambda_2 &= \begin{pmatrix} 0 & -i & 0 \\
i & 0 & 0 \\
0 & 0 & 0 \end{pmatrix}; \\
\lambda_3 &= \begin{pmatrix} 0 & 0 & 1 \\
0 & -1 & 0 \\
0 & 0 & 0 \end{pmatrix}; & \lambda_4 &= \begin{pmatrix} 0 & 0 & 1 \\
0 & 0 & 0 \\
1 & 0 & 0 \end{pmatrix}; \\
\lambda_5 &= \begin{pmatrix} 0 & 0 & -i \\
0 & 0 & 0 \\
i & 0 & 0 \end{pmatrix}; & \lambda_6 &= \begin{pmatrix} 0 & 0 & 0 \\
0 & 0 & 1 \\
0 & 1 & 0 \end{pmatrix}; \\
\lambda_7 &= \begin{pmatrix} 0 & 0 & 0 \\
n & 0 & -i \\
0 & i & 0 \end{pmatrix}; & \lambda_8 &= \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\
0 & 0 & -2 \\
0 & 1 & 0 \end{pmatrix}.
\end{align*}
\]

Each single-system measurement setting has three possible outcomes. For each pair of measurement settings (for the input and the output systems) we simulate \( N \) measurement outcomes. We choose \( N = 10^6 \) for our reference experiment, yielding a total of \( N_{\text{tot}} = 8^2 \times 10^6 = 6.4 \times 10^7 \) measurement outcomes. We denote the corresponding frequency vector by \( (n_{k_{\ell_1,\ell_2,\ell_3}}^{\text{Ref}})_i \), where \( j_i \) labels the measurement setting on system \( i \) and \( \ell_i \) labels the corresponding measurement outcome. We group together all the indices into a collective index \( k \), such that \( n_{k}^{\text{Ref}} \) denotes the number of times the joint POVM effect \( E_{AB}^{(i)} \) was observed.

The corresponding analysis is depicted in Fig. 7, as the curve labeled “100%.” Thanks to the large number of measurements, the distributions \( h(v) \) and \( \mu(v) \) peak sharply around the true value of \( f_0 \). We now ask, how would these distributions look if fewer measurements had been taken? Instead of simulating new outcomes, which would cause the peak to be displaced and would make a comparison more difficult, we artificially rescale the frequency vector \( n_k^{\text{Ref}} \) by a factor \( \alpha \), i.e., we define \( n_k^\alpha = [\alpha \, n_k^{\text{Ref}}] \), where by \( [x] \) we denote the largest integer less than or equal to \( x \). For instance, we may choose \( \alpha = 0.01 = 1\% \) to represent an experiment in which only \( n \approx \alpha \, N = 10^4 \) measurements per setting were sampled, instead of \( N \). While this rescaling of the frequency vector is artificial, the resulting measurement counts are still representative of possible outcomes that one could have sampled if we had simulated directly only \( \alpha \, N \) outcomes per setting; crucially, doing so facilitates comparisons between the different settings. The analysis for a selection of values for \( \alpha \) (given as percentages) is presented in Fig. 7. The corresponding peaks are indeed seen to converge toward the true value of \( f_0 \). For each value of \( \alpha \), we calculate the corresponding quantum error bars \( (v_0, \Delta, \gamma) \), and plot them against \( \alpha \) (Fig. 7, inset). The quantum error bars become a better and tighter description of the true state as the number of measurements increase, as expected.

The quantum error bar \( \Delta \) is the one which is most akin to a “standard deviation,” as in the limit \( \gamma \to 0 \) the fit model (9) becomes a Gaussian. We may investigate the precise scaling of \( \Delta \) as a function of the number of measurements.
by plotting the magnitude of this quantum error bar against the number of measurements in a log-log plot (Fig. 8). We indeed observe a scaling close to $1/\sqrt{N}$, where $n \approx aN$ is the number of measurements, as expected from known results in usual quantum tomography. We expect that by improving the measurement settings, for instance by using adaptive measurements, tighter error bars can be achieved with fewer measurements \cite{13,50–52}. This depiction allows us again to appreciate the convergence to the true value of $f_0$.

APPENDIX G: RELATIONS OF TWO METHODS

In this Appendix, we discuss the theoretical connections between the two methods, specifically the relationship between the densities $\mu(v)$ and $h(v)$. We will use basic notions from measure theory which is available in any standard textbook.

Recall that we use the induced measure $d\sigma_{AB}$ on density matrices $D(H_{AB})$ in the bipartite-state sampling and the measure $d\nu(\psi)$ on Choi state $\mathcal{C}(H_{AB})$ in the channel-space method. It is helpful for the reader to refresh the definition of these measures in Appendix A 2. The following result connects these probability measures; its proof is delayed until the end of this Appendix.

Proposition 5. The measure $d\sigma_{AB}$ factors as $d\sigma_{AB}d\nu(\Lambda_{A\rightarrow B})$ in the sense that for all measurable function $g(\sigma_{AB})$,

$$\int d\sigma_{AB}g(\sigma_{AB}) = \int d\sigma_A \int d\nu(\Lambda_{A\rightarrow B})g(d\sigma_A^{1/2}J(\Lambda_{A\rightarrow B})\sigma_A^{1/2}) \quad (G1)$$

$$\int d\sigma_{A} \int d\nu(\Lambda_{A\rightarrow B})g(d\sigma_A^{1/2}J(\Lambda_{A\rightarrow B})\sigma_A^{1/2}) = \int d\sigma_{A} \int d\nu(\Lambda_{A\rightarrow B})g(d\sigma_{A_{1/2}}^{1/2}J(\Lambda_{A\rightarrow B})\sigma_{A_{1/2}}^{1/2}) \quad (G2)$$

where $d\sigma_A$ is the reduced measure of $d\sigma_{AB}$ via partial tracing and $d\nu(\Lambda_{A\rightarrow B})$ is the uniform measure on channel space induced by $dU_{BA}$ and $\sigma_P = \sigma_A^{T}$.

We remark that intuitively this result is clear: the probability measure $d\sigma_{AB}$ can be “conditioned” on different values of $y = tr_P(\sigma_{AB})$ giving rise to conditional probability measures $d\nu(y|\sigma_{AB})$ and these are recognized as $d\nu(\Lambda)$ by unitary invariance. However, the fact that these events which we are conditioning on have measure zero under $d\sigma_{AB}$ makes the proof more complicated.

Proposition 5 tells us that integrating over all bipartite states according to the measure $d\sigma_{AB}$ can be done separately integrating over all possible input states $\sigma_A$ and over all possible channels $\Lambda_{A\rightarrow B}$, by combining them as $\sigma_A^{1/2} \Lambda_{AB} \sigma_A^{1/2}$ where $\Lambda_{AB} = J(\Lambda_{A\rightarrow B})$. Equivalently, this can be done by separately integrating over all possible transposed input states $\sigma_A$ and over all possible channels $\Lambda_{A\rightarrow B}$ as in (G2). We can use this intuition to relate the two methods presented above.

In order to connect both quantities, we consider the situation depicted in Fig. 9. Assume that for each repetition $j = 1 \ldots n$ the input $\rho_A^j$ is chosen by a measurement on the pure state $|\psi\rangle_A = \sigma_A^{1/2}|\Phi\rangle$ for some given state $\sigma_A$, and the outcome POVM effect $E_P$ was observed. The measurement on the output state of the channel is chosen from some collection of measurements acting only on system $B$ only. Assuming that
the outcome POVM effect $E^I_B$ was observed, the data set $E$ consists of the pairs $(E^I_P, E^I_B)$ for all $n$ repetitions.

Viewing this scenario as an ancilla-assisted scheme (by moving the measurement on $P$ to the end), we can employ the bipartite-state sampling method and calculate $\mu(v)$ by integrating our test function $\delta(f(\rho_{AB}) - v)$ over the full bipartite-state space according to (B15) and (B5):

$$\mu(v) = c_E^{-1} \int d\rho_{AB} L_1(\rho_{AB}|E) \delta(f(\rho_{AB}) - v). \quad \text{(G3)}$$

where

$$L_1(\rho_{AB}|E) = \text{tr}(\sigma_{AB}^{\text{out}} E), \quad \text{(G4)}$$

for $E = \otimes_{j=1}^n E^I_P \otimes E^I_B$.

On the other hand, we can also view this as a prepare-and-measure scheme and use the channel-space method to compute the histogram by (C3) and (C42) as an integration over the space of all channels only:

$$h(v) = c_E^{-1} \int d\rho_{A\rightarrow B} L_2(\rho_{A\rightarrow B}|E) \delta(f_{\text{chann}}(\Lambda) - v). \quad \text{(G5)}$$

where

$$L_2(\Lambda|E) = \prod_{j=1}^n \text{tr}(\Lambda_{A\rightarrow B}(\rho_A^j) E^I_B). \quad \text{(G6)}$$

We may rewrite each factor term using the Choi-Jamiolkowski state of the channel as

$$\text{tr}(\Lambda_{A\rightarrow B}(\rho_A^j) E^I_B) = \text{tr}(\sigma_{PB}^{1/2} \Lambda_{PB} \sigma_{PB}^{1/2} (E^I_B \otimes E^I_j)), \quad \text{(G7)}$$

where $\sigma_P = \sigma_A^\dagger$ and thus

$$L_2(\Lambda|E; \sigma_A) = \text{tr}((\sigma_{PB}^{1/2} \Lambda_{PB} \sigma_{PB}^{1/2}) \rho_A^{\text{out}} E), \quad \text{(G8)}$$

now defining the same operator $E = \otimes_{j=1}^n E^I_P \otimes E^I_B$ as before and where $\sigma_A$ is fixed.

The similarity of (G3) and (G4) with (G5) and (G8) is now more evident. It is worth giving a precise interpretation to both $L_1(\sigma_{AB}|E)$ and $L_2(\Lambda|E; \sigma_A)$. The function $L_1(\sigma_{AB}|E)$ is a probability density on the bipartite-state space with respect to $d\sigma_{AB}$, describing the Bayesian posterior distribution after observing data $E$ for an agent using the uniform prior $d\sigma_{AB}$ (and thus ignoring any prior information about what the input state actually is). On the other hand, $L_2(\Lambda|E; \sigma_A)$ is the posterior distribution in the space of all channels, after observing data $E$ for an agent which is using the prior $d\rho_{A\rightarrow B}$. Yet, Proposition 5 tells us that the prior $d\rho(\Lambda_{A\rightarrow B})$ is precisely the same as the prior in the bipartite-state space corresponding to knowing with certainty that the input state is exactly $\sigma_A$. Indeed, $d\rho(\Lambda_{A\rightarrow B})$ is precisely the measure induced by $d\sigma_{AB} \delta(\text{tr}_B(\sigma'_{AB}) - \sigma_A) / \Lambda_{A\rightarrow B}$.

Hence, the difference between the input-state sampling method and the channel-space method, at least in the current scenario, is exactly the prior information about the input state. In the former, nothing is assumed about the input state other than what can be inferred directly from the measurement data; in the latter, the exact input state is assumed with certainty as represented by the first Dirac delta function in (G9).

Finally, we will prove the following result, which is easily seen to imply the Proposition 5.

Proposition 6. There exists an essentially unique family of probability measures $d\nu(\sigma_{AB})$ on $D(\mathcal{H}_{AB})$ indexed by full rank $y \in D(\mathcal{H}_A)$ such that

$$\int d\sigma_{AB} g(\sigma_{AB}) = \int d\nu(\sigma_A) \int_{\mathcal{H}_B(\sigma_A)} d\nu_B(\sigma_B) g(\sigma_{AB}). \quad \text{(G10)}$$

where $d\sigma_A(y)$ is the reduced measure of $d\sigma_{AB}$ and $\nu_B(\sigma_B)$ denotes the preimage of $y$ under partial tracing $B$. Moreover, each member $d\nu(\sigma_{AB})$ of the family is supported on $\mathcal{H}_B(\sigma_A)$ and actually isomorphic to $d\nu(\Psi)$ on $\mathcal{H}_A \rightarrow \mathcal{H}_B$. These isomorphisms are given by

$$J_1^{-1} : \mathcal{H}_B(\sigma_A) \rightarrow \mathcal{H}_A \rightarrow \mathcal{H}_B \quad \text{(G11)}$$

$$\sigma_{AB} \mapsto J_1^{-1}(d_1^{-1} y^{-1/2} \sigma_{AB})^{-1/2}. \quad \text{(G12)}$$

Proof. Again, it will be convenient to work in the purified picture. By definition, $d\sigma_{AB}$ originates from the uniform spherical measure $d|\phi\rangle_{\mathcal{A}\mathcal{B}}$ induced by the Haar measure $dU_{\mathcal{A}\mathcal{B}}$ by the relation $|\phi\rangle_{\mathcal{A}\mathcal{B}} = U_{\mathcal{A}\mathcal{B}} |\psi_0\rangle$. On the other hand, $d\nu(\Psi)$ comes from the Haar measure $dU_{\mathcal{B}\mathcal{A}}$ via the relation $|\psi\rangle = U_{\mathcal{B}\mathcal{A}} |\psi_0\rangle$.

Consider the partial trace $tr_{\mathcal{B}A} : \text{End}(\mathcal{H}_{\mathcal{A}\mathcal{B}}) \rightarrow \text{End}(\mathcal{H}_A)$. Two things happen under this mapping. First, the measure $d|\phi\rangle_{\mathcal{A}\mathcal{B}}$ admits a push forward along $tr_{\mathcal{B}A}$ denoted as $d\sigma_A(y)$ living on space $D(\mathcal{H}_A)$. Note that this measure $d\sigma_A(y)$ no longer coincides with the Haar-induced (or Hilbert-Schmidt-induced) uniform measure on $D(\mathcal{H}_A)$ (since such measure arises uniquely from the Haar measure $dU_{\mathcal{A}\mathcal{A}}$ acting on $\mathcal{H}_{\mathcal{A}\mathcal{A}}$). Second, the space $\text{End}(\mathcal{H}_{\mathcal{A}\mathcal{B}})$ is partitioned into fibers $tr_{\mathcal{B}A}^{-1}(y)$ over $y \in \text{End}(\mathcal{H}_A)$. Observe that one of such fibers corresponds to the set of purified Choi states $\mathcal{P}C$: take $y = 1/\Lambda_A$. Moreover, if $y \in \text{End}(\mathcal{H}_A)$ is full rank, then the fiber over $y$ is isomorphic to $\mathcal{P}C$. Indeed, the bijection is given by

$$J_1^{-1} : tr_{\mathcal{B}A}^{-1}(y) \rightarrow \mathcal{P}C, \quad \text{(G13)}$$

$$\phi_{\mathcal{A}\mathcal{B}} \mapsto d_1^{-1} y^{-1/2} \phi_{\mathcal{A}\mathcal{B}} y^{-1/2}. \quad \text{(G14)}$$
Note that partial tracing out $A'B'$ gives Choi-Jamiolkowski isomorphisms identifying $tr_{A'B'}^{-1}(y) \subseteq D(H_{A'B'})$ with the space of all quantum processes:

$$J_y^{-1} : tr_{A'B'}^{-1}(y) \rightarrow \mathcal{C}(H_A \rightarrow H_B).$$

$$\sigma_{AB} \mapsto J^{-1}(d_{\sigma_{AB}}^{-1/2}y^{-1/2}),$\]

where $J^{-1}$ is the standard Choi-Jamiolkowski isomorphism identifying $\mathcal{C}(H_{AB})$ with $\mathcal{C}(H_A \rightarrow H_B)$. We stress again that these are isomorphisms only for full rank $y \in D(H_A)$.

The probability measure $d\phi_{AB'}$ then disintegrates into a family of conditional probability measures $d\nu_{S}(\phi_{AB'})$ on each fiber (or preimage over $y$) $tr_{A'B'}^{-1}(y)$ such that

$$\int d\phi_{AB'} g(\phi_{AB'})$$

by change of variables

$$\int d\sigma_A(y) \int_{tr_{A'B'}^{-1}(y)} d\nu_{S}(\phi_{AB'}) g(\phi_{AB'})$$

for all functions $g(\phi_{AB'})$. Moreover, the family $\{d\nu_{S}(\phi_{AB'}) : y \in D(H_A)\}$ is $\sigma_A(y)$ almost everywhere unique and each member $d\nu_{S}(\phi_{AB'})$ is supported on $tr_{A'B'}^{-1}(y)$.

Without loss of generality, we only pay attention to full rank $y \in D(H_A)$ because the set of rank-deficient density matrices $y$ has measure zero under $d\phi_{AB'}$. Here, each fiber $tr_{A'B'}^{-1}(y)$ has been identified with the space $\mathcal{P}^{\mathcal{P}}$ of density matrices. Under this identification, we will show that $d\nu_{S}(\phi_{AB'})$ is almost everywhere equivalent to the uniform measure on channel space $d\nu(\Psi)$. This follows from unitary invariance of $d\nu_{S}(\phi_{AB'})$ and the uniqueness of the Haar measure $d\sigma_{AB'}$. Specifically, since $d(U_{A'B'}^{1}\phi_{AB'}U_{A'B'}) = d\phi_{AB'}$ for all $U_{A'B'}$, we have

$$\int d\phi_{AB'} g(\phi_{AB'}) = \int d(U_{A'B'}^{1}\phi_{AB'}U_{A'B'}) g(\phi_{AB'})$$

where the last equality follows from the fact that the fiber $tr_{A'B'}^{-1}(y)$ is invariant under all $U_{A'B'}$. By uniqueness of the family, we have

$$d\nu_{S}(U_{A'B'}^{1}\phi_{AB'}U_{A'B'}) = d\nu_{S}(\phi_{AB'})$$

This says that each member $d\nu_{S}(\phi_{AB'})$ of the disintegration family is unitary invariant. Due to the uniqueness of the normalized Haar measure, we conclude $d\frac{1}{tr_{A'B'}}(\phi_{AB'}) = d\nu(\Psi)$. In fact, we obtain correspondences between the objects

$$\nu^{tr_{A'B'}} \leftrightarrow \nu^{\mathcal{P}^{\mathcal{P}}},$$

$$\frac{1}{tr_{A'B'}}(\phi_{AB'}) \leftrightarrow \nu(\Psi)$$

induced by $J_y^{-1}$.

Taking partial trace of system $A'B'$ yields the statement of the proposition and completes the proof. \hfill$\blacksquare$

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