I. INTRODUCTION

Entanglement describes quantum correlations with no classical analog, and it underpins many advantages of quantum devices over classical computation, communication, and sensing [1–3], while also being central in many physical phenomena such as phase transitions [4]. However, entanglement is difficult to characterize both theoretically and, even more so, experimentally. The most direct way requires performing quantum state tomography (QST) [5,6] to obtain the density operator $\rho$ describing the state and then use one of several metrics of entanglement that have been proposed. Unfortunately, QST requires a number of measurements that scale exponentially with the qubit number $n$, such that for large $n$ QST becomes intractable. Even for small systems, errors in state preparation and measurement (SPAM) compound the difficulty in identifying $\rho$ with a high accuracy [7–10].

When the goal is more simply to detect whether or not entanglement is present, an attractive alternative is to use so-called entanglement witnesses $W$ [11,12]. In contrast to QST, the resources needed to measure an entanglement witness typically scale more favorably with respect to $n$. The witness operator $W$ is designed to ‘witness,’ that is, detect a specific entangled state $|\psi\rangle$: its expectation value $\langle W_{\psi} \rangle = \text{tr}(\rho W_{\psi})$ is negative for some entangled states, while it is positive for all separable states. While designing a witness for an arbitrary entangled state is difficult, since this would solve the separability problem [13–16], for NPT entangled states $|\psi\rangle$ (states that have negative eigenvalues under positive partial transpose) [17], such as the well-known Bell, GHZ, $W$, and Dicke states, the witness is based on the state fidelity, $\langle F_{\psi} \rangle \equiv \langle \psi|\rho|\psi\rangle$:

$$W_{\psi} = \alpha \mathbb{1} - |\psi\rangle\langle\psi|.$$  

(1)

Here $\alpha$ is the squared maximum overlap of $|\psi\rangle$ with all possible separable states [18] so as to conservatively detect entanglement. This improved experimental feasibility—requiring only the measurement of state fidelity—has led to successful demonstrations of entanglement detection across many platforms [18–24] as well as theoretical improvements and modification of the witness to further improve experimental feasibility [25].

However, one immediate drawback of such ‘state’ witnesses $W_{\psi}$ is that they can severely underestimate the amount of entanglement actually present in $\rho$. While errors in state preparation can indeed lower the entanglement from the target state $\rho = |\psi\rangle\langle\psi|$, some common errors such as local and unitary errors will not change the amount of entanglement. However, the typical witnesses $W_{\psi}$ of the form in Eq. (1) will not capture this entanglement.

Here we use a solid-state two-qubit system to investigate the advantages and limits of entanglement witnesses in the presence of SPAM errors. To achieve more robust entanglement detection, we introduce a new metric, which we call the ‘subspace’ witness $W_s$, that can capture a larger share of entangled states generated in the presence of unitary, local errors. We compare state $\langle W_{\psi} \rangle$ and subspace $\langle W_s \rangle$ witnesses.
observing an improvement in the detection of entanglement by \( W_{\psi} \), which can even provide a stricter bound on entanglement quantification. As an extension, we discuss the subspace witness measurement for specific genuine multipartite entangled (GME) states, such as GHZ, W, and Dicke states, that are compatible with state-fidelity witnesses. We find that quite broadly the subspace witness allows identification of all the many-body coherences, which are often of interest in practical applications such as quantum sensing.

II. WITNESSING TWO-QUBIT ENTANGLEMENT

A. State witnesses \( W_\psi \)

For a two-qubit system, there are four canonical maximally entangled states, the Bell states [2,26,27]. The Bell states \(|\Phi^\pm\rangle, |\Psi^\pm\rangle\) form an orthogonal basis, thus any state (and, in particular, entangled states) can be written in terms of their superpositions. Choosing the computational basis to describe the energy eigenbasis, we can explicitly write a Bell state as \(|\Phi^\pm\rangle = (|k\rangle \pm |\bar{k}\rangle)/\sqrt{2} \), with \(|k\rangle = |00\rangle\,(|01\rangle)\) and \(|\bar{k}\rangle\) the corresponding spin-flipped states. Each pair of Bell states, \(|\Phi^\pm\rangle\) and \(|\Psi^\pm\rangle\), spans a subspace with constant energy. For many applications, such as entanglement-enhanced quantum sensing [3,28,29] or decoherence-protected subspaces [30–35], states inside this subspace are equally beneficial. In particular, we can identify the family of maximally entangled states inside such subspaces, parametrized by a phase \( \varphi \),

\[
|\Phi(\varphi)\rangle = \cos(\varphi/2)|\Phi^+\rangle + i \sin(\varphi/2)|\Phi^-\rangle,
\]

and similarly for \(|\Psi^\pm\rangle\). Here \( \varphi \) describes the phase degree of freedom that leaves unchanged the state desired properties (e.g., for enhanced sensing or decoherence-protected memory, respectively).

Fixing \( \varphi \), we can build a canonical state witness as in Eq. (1) (with \( \alpha = 1/2 \)). This is a good witness to detect the presence of two-qubit entanglement in any state \( \rho \), given that all two-qubit entangled states are NPT. The expectation value of the witness depends on the state fidelity, \( W_\psi = 1/2 - F_\psi \), where the state fidelity is a function of \( \varphi \):

\[
F_\psi = \langle \psi(\varphi) | \rho | \psi(\varphi) \rangle = P + C(\varphi) \leq P + |\rho_{\bar{k}k}|.
\]

Here \( P = 1/2(\rho_{kk} + \rho_{\bar{k}\bar{k}}) \) is the sum of populations in the \(|k\rangle, |\bar{k}\rangle \) subspace and

\[
C(\varphi) = \text{Re}(\rho_{\bar{k}k}) \cos(\varphi) + \text{Im}(\rho_{\bar{k}k}) \sin(\varphi)
\]

are the related coherences, with tan\( (\theta_{\bar{k}k}) = \text{Im}(\rho_{\bar{k}k})/\text{Re}(\rho_{\bar{k}k}) \). The coherence \( C(\varphi) \) is maximum only for \( \theta_{\bar{k}k} = -\varphi \). Unfortunately, \( \theta_{\bar{k}k} \) might be unknown due to the unitary component of SPAM errors. Then, while \( P < 1/2 \) always reflects a suboptimal (or absent) entanglement, \( C(\varphi) \) might even be negative, although the state is maximally entangled. This leads to an underestimate not only of the entanglement unless \( \theta_{\bar{k}k} = -\varphi \) but, more critically, of the useful entanglement for many quantum tasks, as often the exact value of \( \theta_{\bar{k}k} \) is unimportant.

B. Subspace witnesses \( W_i \)

As a way to improve upon entanglement detection by state witnesses, we propose a subspace witness measurement that becomes insensitive to some unitary SPAM errors:

\[
\langle W_i \rangle = \min\langle W_\psi \rangle = \alpha - \max \langle \psi(\varphi) | \rho | \psi(\varphi) \rangle.
\]

We call this a subspace witness as, for any state in the subspace spanned by the relevant entangled-state basis, the witness is able to detect whether or not it is entangled. The subspace witness can thus be considered an intermediate metric between state witnesses and entanglement measures: while entanglement measures provide a quantitative estimate of the entanglement amount, typically by optimizing over all possible local unitaries, the subspace witness optimizing over a set of local unitaries that are of relevance for a particular quantum information task, thus detecting entanglement more robustly than the state witness, while still maintaining an efficient protocol.

Indeed, to experimentally obtain the subspace witness by maximizing the fidelity over the subspace, one needs to simply perform multiple state witness (or fidelity) measurements. That is, improved entanglement detection comes at the cost of additional measurements. Still, the number of measurements is much smaller than for QST. For two-qubit entanglement, Eqs. (3) and (4) show that there are three unknowns: \( \text{Re}(\rho_{\bar{k}k}), \text{Im}(\rho_{\bar{k}k}), \) and \( P \). Thus, three measurements, e.g., at \( \varphi = \{0, \pi/2, \pi\} \), fully identify \( P \) and \( |\rho_{\bar{k}k}| \), thus yielding the subspace-optimized entanglement witness. While the advantage for two qubits is not large, it quickly becomes substantial for larger systems, as we see in Sec. IV.

Having discussed the idea of the subspace witness we now describe our experimental system and the experimental protocol to measure \( W_i \).

III. EXPERIMENTAL GENERATION AND DETECTION OF ENTANGLEMENT

A. Entanglement generation

We generated entanglement in a solid-state two-qubit system, comprising two electronic spin impurities in diamond. The qubits are given by two electronic-spin levels of a single nitrogen-vacancy (NV) center, weakly interacting with a nearby, optically dark, electronic spin-1/2 defect in diamond. This spin qubit has been earlier characterized as an electronic-nuclear spin defect [36]. Here we neglect the nuclear spin degrees of freedom and remove the electronic spin dependence on their state by applying two-tone microwave pulses detuned by the nuclear hyperfine strength.

To experimentally prepare an entangled state \( \rho \), we follow the protocol described below whose details are provided in [29,37]. Starting from the thermal equilibrium state of the two spins, we first initialize the two-qubit system to \( \rho_0 \) via a Maxwell demon-type cooling scheme. Specifically, we first initialize the NV spin to a high-purity state using spin-nonconserving optical transitions under laser illumination [38,39]. Then we swap its state with the dark electronic spin \( X \) and reinitialize the NV . After the initialization step \( I \), we apply an entangling operation \( U \) to prepare the desired state \( \rho \) [Fig. 1(a)].

Both the SWAP in \( I \) and the entangling gate in \( U \) are achieved by Hartmann-Hahn cross-polarization (HHCP), which exploits the spin coupling described by
by decoherence, but also by local unitary rotations. Then the prepared state $\rho$ might not be as desired, and a state witness might underestimate the entanglement present. To compound these issues, as we explain in the next section, similar control operations are needed to measure entanglement, given the available observables. Thus to partially relieve these SPAM errors, we show that it is beneficial to use subspace witnesses.

In the following, we measure both the state ($W_\psi$) and the subspace ($W_s$) witnesses and observe the presence of unitary, local errors, thus motivating the use of ($W_\psi$).

### B. Witness measurements

Here we discuss the experimental protocol for measuring both the state witness ($W_\psi$) and the subspace witness ($W_s$) in our system. Ideally, to measure the Bell-state fidelity $F_\psi$ that enters in ($W_\psi$) we would want to measure the observable $M'_s = |\Phi^+\rangle\langle\Phi^+|$ such that the experimental signal would directly yield the Bell fidelity: $S = \text{tr}(M'_s \rho U) = (\Phi^+|\rho|\Phi^+)$. While unfortunately this is typically not possible, we can obtain $M'_s$ by a suitable unitary transformation of any joint projective measurement operator on the two qubits [45]. For example, consider an experimental system with the joint projective measurement operator $M_{0\rightarrow 2} = |0\rangle\langle0|^{0\rightarrow 2}$. Prior to measurement, we evolve the state of interest $\rho$ under a unitary disentangling gate $U^\dagger$, such as $U = C_1\text{NOT}_2 \cdot H_1$ (where $H_1$ is the Hadamard gate applied to qubit i). This is equivalent to transforming the bare measurement operator $M_2$ into $M'_2 = U M_2 U^\dagger = |\Phi^+\rangle\langle\Phi^+|$, as desired.

Unfortunately, our hybrid system lacks a joint projective measurement operator such as $M_2$, as we can only directly measure the state of the NV center, $M = |0\rangle\langle0| \otimes I$. In this case, even with universal control on the two-qubit system, we cannot measure the Bell fidelity in a single measurement. Therefore, here we introduce a protocol to reconstruct the Bell-state fidelity that exploits measuring three correlators, $\langle \sigma_{\alpha}^a \sigma_{\alpha}^b \rangle$, with $\alpha = \{x, y, z\}$.

In the experiments, we measure $M' = U M U^\dagger = \sigma_1^a \sigma_2^a$, with $U = C_1\text{NOT}_2$ and the bare operator $M = \sigma_1^a = |0\rangle\langle0| - |1\rangle\langle1|$ obtained from the difference signal of measuring the NV states 0 and $-1$. The controlled-\text{NOT} gate is realized by exploiting evolution under the two-spin interaction Hamiltonian $H_{\text{int}}$, via the pulse sequence $R_\alpha(\frac{\pi}{4}) \cdot e^{-iHt} = R_\alpha(\theta)$, with $t = \pi/4d$ and $R_\alpha(\theta)$ collective rotations of the two spins along the axis $\alpha$ for an angle $\theta$. The other two correlators $M' = \sigma_1^a \sigma_2^a$ for $\alpha = \{x, y\}$ can be measured by adding a global $R(\pi/2)$ rotation along the $\{y, z\}$ directions, respectively.

Given a joint projective measurement as discussed above, globally rotating the disentangling gates along $z$, $M'_2(\psi) = |\Phi^+\rangle\langle\Phi^+| R_z(\psi) |\Phi^+\rangle\langle\Phi^+| R'_z(\psi)$, would yield the required measurements to reconstruct the subspace witness. Indeed, recall that to measure $W_s$ from the subspace-maximized Bell-state fidelity $\text{max}_\psi [F_\psi(\psi)]$ requires multiple fidelity measurements in order to learn the magnitude of the coherence, $|\rho_{jk}|$. Similarly, if only one of the two qubits is observable, $M_{12}^\alpha$, one could use the correlator $\langle \sigma_1^a \sigma_2^a \rangle$ and their $R_\alpha(\psi)$ rotations to extract $W_i$.

First, we note that $\sigma_1^a \sigma_2^a$ is invariant under $R_y$, and indeed it yields the population $P = (1 + |\langle \sigma_1^a \sigma_2^a \rangle|)/4$, which is independent of $\psi$. To extract $C(\psi)$ here we used the same HHCP-based disentangling gate $U$ that generated the entanglement, resulting in the following signal $S = \langle M' \rangle$ up to a
constant:

\[
S(\phi) = \text{tr}(\rho U(\phi) MU^{-1}(\phi))
\]
\[
= \frac{1}{2} |\sigma^x_1| \cos(\phi) - \frac{1}{2} |\sigma^y_1| \sin(\phi)
\]
\[
= 2 \text{Re}(\rho_{14}) \cos(\phi) - 2 \text{Im}(\rho_{14}) \sin(\phi)
\]
\[
= 2(C(\rho)). \tag{6}
\]

Here \(\rho_{14} = \langle 00|\rho|11\rangle\) is the coherence of interest, and the (undesired) constant offset under HHCP is given by \(|\sigma^x_1 - \sigma^z_1|/2\), thus yielding a total of four measurements required to determine \(W_s\) (while only three would be necessary if a joint projective measurement were available; see Appendix A for experimental details and signal derivation.)

C. Experimental results

We now discuss the measurement results when attempting to create the Bell state \(|\Phi^+\rangle\), which results in the generation of the state \(\rho\). First, we measure the state witness \(\langle W_s(\phi = 0)\rangle\), which requires the measurement of the Bell-state fidelity \(\langle \Phi^+|\rho|\Phi^+\rangle = \frac{1}{4}(1 + |\sigma^x_1|^2 + |\sigma^y_1|^2 - |\sigma^z_1|^2)).\) From the three measurements shown in Fig. 2(a), we obtain \(F_{\Phi} = 1/4(1 + 0.497 + 0.2142 + 0.5857) = 0.57421(4) > 1/2\). From this fidelity we have \(\langle W_s\rangle = 0.07,\) which successfully detects entanglement.

Still, this measurement might underestimate the amount of entanglement, due to coherent, local errors. We thus measure the subspace witness \(\langle W_s\rangle\), to extract the coherence \(|\rho_{\text{coherent}}| = |\langle 00|\rho|11\rangle|\) [see Fig. 2(b) and Appendix A for experimental details and signal derivation]. We obtain a maximum fidelity of \(\langle \Phi|\rho|\Phi\rangle = 0.6827(4),\) corresponding to \(W_s = -0.1827(4),\) having optimized over the \(|\Phi\rangle\) subspace. This indicates that we indeed have coherent errors affecting our state preparation and measurement process.

The subspace witness can also help determining the entanglement coherence time. The (entangled) prepared state \(\rho(\tau = 0)\) will evolve under the environment influence and the natural (or engineered) Hamiltonian. By measuring the subspace witness after a time \(\tau\), one can then detect whether entanglement is still present, at the net of local, unitary evolution. By setting \(|\rho_{\text{coherent}}(\tau^*)| = \alpha\) we can further define a threshold time \(\tau^*\) after which entanglement in \(\rho(\tau \geq \tau^*)\) is no longer witnessed.

To simplify the measurement of the subspace witness, we can apply the phase \(\phi\) rotation at each time point measured, such that \(\phi = \nu \tau\). Then, from the decay of the oscillations in \(C\), one can directly extract \(W_s\) at \(\tau = 0\) and the characteristic time \(\tau^*\). We note that in general both coherences \(C\) and population \(P\) will decay for open quantum systems, but only extraction of the coherences at various \(\nu\) is needed to reconstruct \(W_s\). In our experiments [see Fig. 2(c)], we thus measure the phase-modulated decay of the coherence. As the main decoherence source is dephasing, which leaves populations intact, we simply verify that \(P(\tau)\) is constant.

In experiments, we compare \(P(\tau) = (1 + |\sigma^x_1|^2)\) at \(\tau = 0\) and \(\tau = 40 \mu s\) and, as expected, observe no decay in \(|\sigma^z_1|^2\). We then study the entanglement decay under a spin echo [29,46] of varying duration \(\tau\). The prepared entangled state \(\rho\) evolves under \(H_{\text{int}}\) (which would not affect the ideal state \(|\Phi\rangle\)) and decoherence. We then apply the disentangling gate (measurement) modulated at \(\phi = \nu \tau\) with \(\nu = 15 \text{kHz}\). A simple Gaussian decay fit yields a characteristic decoherence time \(T_2 = 31(3) \mu s\) of the double-quantum coherence \(|\langle 00|\rho|11\rangle|\), such that taking a constant \(P = 0.374(1),\)
Finally, we note that the subspace witness provides a stricter bound on the amount of entanglement generated. While this might seem intuitive, we remark that a ‘general’ way to relate entanglement detection to existing quantifiers is not known. However, for the two-qubit case, it has been shown [2] that one can relate Bell fidelities $F_\psi$ to the entanglement of formation and, thereby, to any other related metrics such as the well-known concurrence. More specifically, the lower bound to the two-qubit concurrence is $C_2(\rho) \geq \max(0, -2W_\psi)$, with $C_2 = 1$ for maximally entangled states. This relation makes it clear that to obtain an entanglement measure one should optimize over all state witnesses. While the subspace witness only optimizes over a restricted set of states, it still provides a stricter bound than the state witness, $C_2 \geq -2W_\psi$.

IV. EXTENSION TO SPECIFIC GENUINE MULTIPARTITE ENTANGLED STATES

We wish to extend the notion of subspace witness to those multipartite entangled states that allow entanglement detection by witnesses $W_\psi$. To this end, we first parametrize a multipartite entangled state $|\psi(\vec{\varphi})\rangle$ in the computational basis $|k\rangle$ as

$$|\psi(\vec{\varphi})\rangle = \sum_{k=0}^{d-1} a_k e^{-i\psi_k}|\vec{k}\rangle,$$

where $a_k$ is the probability amplitude, $\vec{\varphi}_k = \vec{k} \cdot \vec{\varphi}$ is the $\{k\}$-dependent phase given a preset $n$-length vector of phases $\vec{\varphi} = \{\varphi_k\}$ for every $j$th qubit, and $d$ is the dimension of the subspace spanned by the set $\{|k\rangle\}$ specifying $|\psi\rangle$. While such an expression could describe any pure state of $n$ qubits, we are interested in NPT entangled states $|\psi\rangle$ for which state witnesses $W_\psi$ are valid, such as GHZ, $W$, or Dicke states. For instance, for a general $n$-qubit GHZ state $|\psi_\lambda\rangle = (|k\rangle + e^{-i\psi}|\bar{k}\rangle)/\sqrt{2}$, we have $a_k = 1/\sqrt{2}$ and the subspace of interest is spanned by $\{|k\rangle, |\bar{k}\rangle\}$ of dimension $d = 2$. For a $W$ state $|W(n)\rangle = \sum_{k=1}^{n} e^{-i\psi_k}|k\rangle/\sqrt{n}$, we have $a_k = 1/\sqrt{n}$ with the subspace of dimension $d = n$ spanned by states $|\{k\}\rangle$ in the one-excitation manifold.

Given this parametrization, the fidelity $F_\psi$ again reduces to a simple expression, $F_\psi(\vec{\varphi}) = P + C(\vec{\varphi})$, with $\vec{\varphi} = \{\varphi_k\}$ for $k = 1, \ldots, d$. Here, $P = \sum_{k=0}^{d-1} a_k^2$ and $C(\vec{\varphi})$ is just a sum similar to Eq. (4) extended to all many-body coherences $\rho_{jk}$ of interest,

$$C(\vec{\varphi}) = 2 \sum_{j=0}^{d-1} \sum_{k>j}^{d} a_j a_k \bar{c}_{jk} \cdot \bar{c}_{kj},$$

where $\bar{c}_{jk} = (\text{Re}(\rho_{jk}), \text{Im}(\rho_{jk}))$ and $\bar{c}_{kj} = (\cos \varphi_{kj}, \sin \varphi_{kj})$ with $\varphi_{kj} = \varphi_k - \varphi_j$.

As in the $n=2$-qubit case, to extract the ‘subspace’ witness $\{W_j\}$, we must identify all $d(d-1)/2$ many-body coherences $\rho_{jk}$ by again solving the set of linear equations given by multiple measurements $F_\psi(\vec{\varphi})$, containing a total of $d(d-1) + 1$ unknowns.

Therefore the subspace dimension $d$, and its scaling with the number of qubits $n$, sets the limit of which entangled states we can obtain the time $\tau^* = 33(3) \mu s$ until which entanglement can be detected.

While we have shown that we can create entanglement in our system, the state fidelity is not optimal. To improve the quality of entanglement in our hybrid system and investigate the source of nonideality, we probe the entanglement as a function of the repetitive state initialization steps (Fig. 3). In this way, we can distinguish between initialization $I$ and control errors in $U$. We repeat the HHCP plus laser polarization block $N = \{1, 3, 5\}$ times to create $\rho_0(N)$, then prepare the entangled state $\rho_0(N)$ and measure the witness $W_\psi(N)$ with fixed control operations. With increasing $N$ we observe an increase in the overall fidelity $F_\psi(N)$, due to an increase in both the population difference [inferred from $P(N)$] and the double-quantum coherence $C(N)$. We note that explicitly verifying an increase in coherence is of practical importance, because specific applications, e.g., classical field sensing with GHZ states, depend strictly on the amount of coherence (not necessarily fidelity) of the entangled state.
we can tackle, given that we want to be efficient with respect to QST.

For GHZ states, \(d = 2\) is constant and independent of \(n\); in other words, for any \(n\)-qubit quantum processor, one can extract the subspace-maximized GHZ witness with just three measurements, a very efficient protocol. Indeed, experimentally such subspace-optimized fidelity has been observed with \(\sim 20\) qubits in superconducting and neutral atom systems [21–23]. We also note that using a 10-qubit register in diamond up to a 7-qubit GHZ state was witnessed with state fidelity 0.589(5) [24], which could be further optimized with such subspace witness measurement.

For W states, \(d = n\): since this is still polynomial in \(n\), all subspace-optimized \(W\) witnesses will be efficient with respect to QST. In contrast, for Dicke states, \(d = \binom{n}{k} = \binom{n}{k}^{\binom{n}{k}}\) with \(k\) excitations: therefore, only for very lowly (highly) excited Dicke states will \(W\) prove efficient with respect to QST.

While for concreteness we give examples with these well-known GME states, they also highlight the motivation behind improving entanglement detection via subspace witnesses (as alluded to above with the two distinct Bell subspaces): namely, that the amount of useful entanglement for the specific QIP task, i.e., considering entanglement as a resource, often depends on the total magnitude of the quantum coherence within that subspace, independent of relative phases that could result from unknown local SPAM errors. This idea is reflected in entanglement quantifiers, which is independent of local unitary formations, but more concretely seen in well-known QIP tasks of interest. For instance, in entanglement-enhanced sensing of classical fields, it can be shown that the sensitivity depends on the magnitude of the many-body coherence in the GHZ subspace or, for improved quantum memories via a decoherence-free subspace, e.g., with \(W\)-like states, the lifetime of quantum information depends on the magnitude of the coherences. Thus, the subspace witness, while alleviating the problem of local unitary control errors present during entanglement detection measurements, better conveys the notion of entanglement as a resource, almost as an intermediate between entanglement detectors and quantifiers.

In fact, this link could be made more concrete by noting that the subspace witness also facilitates quantification of the lower bound of entanglement, with additional measurements of population terms. More specifically, the lower bound to entanglement, with additional measurements intermediate between entanglement detectors and quantifiers, conveys the notion of entanglement as a resource, almost as a larger amount of witness violation than does \(W\) at the cost of additional measurements, while being efficient with respect to state tomography. Conceptually, the appropriate subspace is chosen for relevance to particular quantum information tasks, and because the subspace witness yields a value optimized over local unitaries within the subspace, it could be viewed as an intermediate metric between an entanglement measure (invariant under local unitaries) and a typical witness (dependent on local unitaries). Experimentally, using a two-qubit solid-state system composed of electronic spins, we observe a significantly improved detection by the subspace witness, motivating the use even for small quantum systems which may have nonnegligible local unitary errors. Finally, we extend the subspace witness to improve detection for specific genuine multipartite entangled states, which may aid near-term NISQ devices to better characterize their performance for applications involving specific entangled states of interest. Because the subspace witness essentially identifies the true (i.e., maximum) many-body coherences of interest, it not only guarantees improved entanglement witness detection but also aids in improving the lower bounds to entanglement via experimentally friendly metrics such as genuine multipartite entangled concurrence or the separability criteria, which require the knowledge of only specific matrix elements of \(\rho\).

V. CONCLUSION

Typical entanglement witnesses based on fidelity measurements to target entangled states \(|\psi\rangle\)—which we call ‘state’ witnesses \((W_\psi)\)—while being efficient detectors of entanglement, can underestimate the actual entanglement present in \(\rho\) due to local, unitary errors in state preparation and measurement. Therefore, we develop the idea of the ‘subspace’ witness \((W_s)\), based on two-qubit systems, which strictly observes a larger amount of witness violation than does \((W_\psi)\) at the cost of additional measurements, while being efficient with respect to state tomography. Conceptually, the appropriate subspace is chosen for relevance to particular quantum information tasks, and because the subspace witness yields a value optimized over local unitaries within the subspace, it could be viewed as an intermediate metric between an entanglement measure (invariant under local unitaries) and a typical witness (dependent on local unitaries). Experimentally, using a two-qubit solid-state system composed of electronic spins, we observe a significantly improved detection by the subspace witness, motivating the use even for small quantum systems which may have nonnegligible local unitary errors. Finally, we extend the subspace witness to improve detection for specific genuine multipartite entangled states, which may aid near-term NISQ devices to better characterize their performance for applications involving specific entangled states of interest. Because the subspace witness essentially identifies the true (i.e., maximum) many-body coherences of interest, it not only guarantees improved entanglement witness detection but also aids in improving the lower bounds to entanglement via experimentally friendly metrics such as genuine multipartite entangled concurrence or the separability criteria, which require the knowledge of only specific matrix elements of \(\rho\).

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APPENDIX A: EXPERIMENTAL DETAILS

AND SIGNAL DERIVATION

1. State witness \((W_\psi)\) measurement

Our experimental system is a hybrid \(n = 2\)-qubit system in which only the first qubit can be directly observed with the (bare) operator \(M = \sigma_1^x \otimes I\), with Pauli operators \(\sigma_{jx,y,z}^{n,2}\) on the \(j\)th qubit. Thus, to measure any two-body correlator \(\langle \sigma_1^x \sigma_2^y \rangle\) (e.g., to reconstruct Bell-state fidelities), one can evolve \(M\) under a two-qubit gate. Given that in our system the qubits interact by \(H_{int} = \delta_0 \sigma_1^z / 2\), a simple experimental sequence involving only single-qubit \(\pi/2\) rotations and free evolution under \(H_{int}\) yields any desired two-qubit correlator.
In other words, first defining a single-qubit rotation on the \( j \)th qubit along axis \( \psi \) by \( R_{\psi,j}(\theta) = e^{-i\theta/2[\sigma_j^x \cos(\psi) + \sigma_j^y \sin(\psi)]} \), a simple pulse, \( U(\psi) = R_{\psi,1}(\pi/2) \cdot e^{-iH_{\text{int}} \cdot \rho/\Phi_1 (\pi/2)} \), will suffice. To mitigate dephasing during free evolution, however, one can also insert a global \( \pi \) pulse(s) on all qubits during the free evolution so as to decouple from the environment bath of spins. Therefore, we insert one \( \pi \) pulse in the middle of the free evolution for both NV and X spin, resulting in \( U(\psi) = R_{\psi,1}(\pi/2) \cdot e^{-iH_{\text{int}}/\Phi_1 \cdot \rho/\Phi_1 (\pi/2)} \cdot R_{\psi,1}(\pi/2) \). Therefore, the effective measurement operator \( M' \) is:

\[
M' = U(\psi)MU(\psi) = \cos(\psi)\left[\sigma_1^x \cos(dt) + \sigma_1^z \sin(dt)\right] + \sin(\psi)\left[\sigma_1^y \cos(dt) - \sigma_1^z \sin(dt)\right].
\]

Thus overlapping \( \rho \) with \( M' \) at the optimal time \( t = (2\pi)/(4d) \) with \( \psi = 0 \) gives the desired signal \( S \):

\[
S = \text{tr}(\rho M') = \langle \sigma_1^z \rangle.
\]

2. Subspace witness (\( W \)) measurement

Measuring the fidelity \( F_{\psi}(\psi) = P + C(\psi) \) similarly involves evolving the bare measurement operator \( M \) under a combination of single-qubit rotations and two-qubit gates. In the experimental work shown, we utilize HHCP to realize the two-qubit gate to both generate and detect entanglement by the subspace witness. More specifically, the evolution under HHCP can be described by the Hamiltonian \( H_{\text{HHCP}} = d(\sigma_1^x \sigma_2^x \pm \sigma_1^y \sigma_2^y)/4 \), in which the sign determines in which subspace (either \( \Phi \) or \( \bar{\Phi} \)) the evolution will occur. For instance, choosing the \( \Phi \) subspace, the bare measurement operator under the pulse sequence \( U(\psi) = R_{\psi,1}(\psi) e^{-iH_{\text{int}}/\Phi_1 R_{\psi,1}(\psi)} \) evolves the desired two-body operators:

\[
2M'(\psi) = 2U(\psi)MU(\psi) = +\left( \sigma_1^x - \sigma_2^x \right) + \left( \sigma_1^y \cos(dt) + \sigma_1^z \sin(dt) \right) - \left( \sigma_1^y \sin(dt) - \sigma_1^z \cos(dt) \right).
\]

Taking the overlap of \( \rho \) with \( M'(\psi) \) at the optimal time \( t = (2\pi)/(4d) \) gives the signal \( S \) carrying the desired information on the two-body coherence:

\[
2S(\psi) = 2\text{tr}(\rho M'(\psi)) = +\langle \sigma_1^x - \sigma_2^x \rangle + \langle \sigma_1^y \sigma_2^y - \sigma_1^y \sigma_1^y \rangle \cos(\psi) - \langle \sigma_1^y \sigma_2^y + \sigma_2^y \sigma_1^y \rangle \sin(\psi) + 4\text{Re}(\langle 00|\rho|11 \rangle) \cos(\psi) + 4\text{Im}(\langle 00|\rho|11 \rangle) \sin(\psi) + \langle \sigma_1^x - \sigma_2^x \rangle + 4C(\psi).
\]

where the first term makes up the (undesired) constant.

APPENDIX B: EXAMPLE FOR BELL STATES

Here we give an example for the two-qubit case (for which an analytical expression of entanglement measure \( C_2 \) can be obtained) that shows that the subspace entanglement witness \( W_j \) detects a larger share of entangled states than does a typical state witness \( W_p \), also with a larger violation.

Consider a generic state in the subspace spanned by one pair of Bell states (e.g., \( |\Phi^\pm\rangle \)). The state can be written as

\[
\rho_\Phi = \frac{1}{2} + \frac{\epsilon}{2} \sin(\phi_0 \cos \theta \Phi_1 + \sin \phi_0 \sin \theta \Phi_2 + \cos \phi_0 \Phi_3),
\]

where \( \Phi_i \) are Pauli matrices in the (sub)space, e.g., \( \Phi_i = |\Phi^i\rangle\langle\Phi^i| + |\bar{\Phi}^i\rangle\langle\bar{\Phi}^i| \), and \( 1_\Phi \) is the identity in the subspace.

\[
\rho_\Phi(\epsilon, \phi_1, \phi_0) \] is uniquely defined in the range \( \epsilon \geq 0, \theta \in [0, 2\pi], \phi_0 \in [0, \pi] \), where \( \epsilon = 0 \) indicates a classical mixture. To gain a bit of insight into this generic state in the chosen subspace, we note that the double-quantum coherence \( C \) (necessarily nonzero for entanglement) is given by

\[
C = \langle 00|\rho_{\Phi}(\epsilon, \phi_1, \phi_0)|11 \rangle = (\epsilon/2) \sin \phi_0 + i \sin \phi_0 \sin \theta.
\]

This shows that maximum entanglement at a given \( \epsilon \)—at \( |C| = \epsilon/2 \)—occurs when either \( \sin \phi_0 = 0 \) or \( \sin^2 \theta = 1 \). The former indicates a state with fully real coherence \( C = \epsilon/2 \), and the latter case refers to a more general complex coherence \( C = e^{i\epsilon/2} \).

Now we calculate the concurrence \( C_2 \) for this state, which reveals that the state \( \rho_\Phi (\epsilon > 0, \phi_1, \phi_0) \) is in general entangled (except at a special point at \( \theta = 0, \phi_0 = \pi/2 \)). More specifically, given \( C_2 = \sqrt{a + b|\epsilon|} - \sqrt{a - b|\epsilon|} \), where

\[
a = (1 - \epsilon^2 \sin^2 \phi_0 (1 + \cos 2\theta))/4,
\]

\[
b = \sqrt{\rho_b b_2}/2,
\]

\[
b_1 = 2 - \sin^2 \phi_0 (1 + \cos 2\theta) \geq 0,
\]

\[
b_2 = 2 - \sin^2 \phi_0 (1 + \cos 2\theta) e^2 \geq 0,
\]

the equality for \( b_1, b_2 = 0 \) is achieved at the mentioned special point at \( \epsilon = 1 \). Since \( C_2 > 0 \), which simplifies to \( 2b|\epsilon| > 0 \), indicates entanglement, we see in general that \( \rho_\Phi (\epsilon > 0, \phi_1, \phi_0) \) is entangled due to positivity of \( b \). Therefore, an ideal entanglement witness should detect all of \( \rho_\Phi (\epsilon > 0, \phi_1, \phi_0) \) (except at the special point).

Now let us first examine the state witness \( W_\psi \) for the ‘range’ of states it can detect as well as its violation. Here the target state is either \( |\psi\rangle = |\Phi^\pm\rangle \); we choose \( |\Phi^+\rangle \), as it makes no difference in the detectable range or the degree of violation. More specifically, we see that

\[
\langle W_\psi \rangle = \frac{1}{2} - \langle \Phi^+|\rho_\Phi|\Phi^+ \rangle = -\frac{\epsilon}{2} \cos \phi_0.
\]

Therefore, assuming \( \epsilon \neq 0 \), the range of detectable states for \( \langle W_\psi \rangle = \phi_0 \in [0, \pi/2] \), while all of \( \phi_0 \in [0, \pi] \) are entangled as seen from concurrence. Furthermore, the state witness, being oblivious to \( \theta \), will underestimate or completely miss all the entanglement in the imaginary part of the coherence.
Finally, we show that the subspace witness $W_s$ captures a larger share of entangled states, namely, all the entangled states in the subspace, also with a larger violation. The subspace witness is

\[
(W_s) = \frac{1}{2} \max_{\phi} (\Phi(\phi)) \rho_\phi |\Phi(\phi)\rangle
= -\frac{\epsilon}{2} \max_{\phi} \{ \sin \theta \sin \varphi_0 \sin \phi + \cos \varphi_0 \cos \phi \}
= -\frac{\epsilon}{2} (\sin^2 \theta \sin^2 \varphi_0 + \cos^2 \varphi_0)^{1/2} \max_{\phi} \{ \cos(\varphi - \varphi') \}
= -\{ |\langle 00 | \rho_\phi (\epsilon, \theta, \varphi_0) |11 \rangle | \max_{\phi} \{ \cos(\varphi - \varphi') \} \}
= -\{ |\langle 00 | \rho_\phi (\epsilon, \theta, \varphi_0) |11 \rangle | \}
\]

(B4)

In other words, the subspace witness detects with maximum violation $|C|$ all of $\rho_\phi (\epsilon > 0, \theta, \varphi_0)$ except at the special (unentangled) point mentioned above. We make one note: the analytical maximization in the third equality is realized at a single $\phi$ since there is only a single coherence $\langle 00 | \rho | 11 \rangle$ to overlap with. Of course, experimentally one does not know the optimal $\varphi$ a priori, so as discussed in the text, such dimension $d = 2$ entangled states require three measurements to measure the subspace witness. An explicit example to measure the subspace witness with multiple coherence terms is discussed in the next section.

**APPENDIX C: EXAMPLE $(W_s)$ MEASUREMENT SCHEME FOR TARGET $|W(n = 4)\rangle$**

Recall that a single-fidelity measurement yields $\langle \psi(\phi) | \psi(\psi) \rangle = P + C(\phi)$, where

\[
C(\bar{\psi}) = 2 \sum_{j=0}^{d-1} \sum_{k=j}^d a_j a_k \bar{\psi}_j \bar{\psi}_k
= 2 \sum_{j=0}^{d-1} \sum_{k=j}^d a_j a_k (\text{Re}(\rho_{jk}) |\psi_{kj}\rangle + \text{Im}(\rho_{jk}) \sin(\varphi_{kj})),
\]

(C1)

with $\psi_{kj} = \psi_k - \psi_j$, where $\varphi_k = \bar{k} \cdot \bar{\psi}$ is the state $|\bar{k}\rangle$-dependent phase given a preset $n$-length vector $\bar{\psi} = \{ \theta_m \}^n_m$, with the phase $\theta_m$ on the $m$th qubit in general. Here, assuming control over the $n$-length vector $\bar{\psi}$ assumes universal control over all qubits such that individual single-qubit control can imprint an arbitrary phase $\theta_m$ along $z$.

Control over individual qubit phases $\theta_m$ allows a simple method to reconstruct the subspace-optimized fidelity and thus improved entanglement detection. For instance, consider the general definition of the $W$-entangled state $|W(n)\rangle = \sum_k e^{-i \phi_k} |k\rangle / \sqrt{n}$. Given single-qubit phase control $\bar{\psi} = \{ \theta_m \}^n_m$, the $W$ state can be written in a more experimentally friendly manner as $|W(n)\rangle = \sum_m e^{-i \phi_m} |m\rangle / \sqrt{n}$, where $m$ indicates excitation on the $m$th qubit. This ‘practical’ definition allows simple parametrization of the fidelity $F(\bar{\psi})$ measurements so as to carefully (with minimal measurements) reconstruct $\max_{\bar{\psi}} [F(\bar{\psi})]$. Given multiple ways to reconstruct $\max_{\bar{\psi}} [F(\bar{\psi})] = P + \max_\phi [C(\bar{\psi})]$, here we discuss one simple approach to reconstruct $\max_{\bar{\psi}} [C(\bar{\psi})]$. First, we note that a judicious choice of $\bar{\psi}$, as seen from Eq. (C1), decouples equations containing either only the real or only the imaginary parts of coherences $\rho_{jk}$. More specifically, $C(\bar{\psi})$ contains either only $\text{Re}(\rho_{jk})$ or only $\text{Im}(\rho_{jk})$ if one measures either $\varphi_0 = \{ 0, \pi \}$ or $\varphi_0 = \{ \pm \pi \}$, respectively. Note that this (arbitrary) restriction to a binary set will accordingly reduce the maximum number of unique fidelity measurements (equations) available to solve for $|\rho_{jk}\rangle$.

More specifically, given $d |k\rangle$ vectors describing the target $|\psi\rangle$, there are now $(d - 1)$ relative phases $e^{-i \theta_m}$, therefore the binary set of inputs allows at most $2^{d-1}$ unique fidelity measurements (equations). Since the number of unknown parameters needed to identify $|\psi\rangle$ is $d(d - 1) + 1$, one must make sure the available equations (in this case $2^{d-1}$) outnumber $d(d - 1) + 1$. The binary restriction of inputs allows this for GHZ, $W$, and a subset of lowly (highly) excited Dicke states. Instead, the subspace dimension $d$ of intermediate excited Dicke states increases nonpolynomially with the qubit number $n$, such that $d(d - 1) + 1 > 2^{d-1}$. In this case, the number of measurements required for entanglement witnessing tends towards that required for state tomography, thus defeating the very purpose of entanglement witnesses.

Finally, we explicitly describe the subspace witness measurement scheme for an $n = 4$-qubit $W$-entangled state:

\[
|\psi(\bar{\psi})\rangle = (10001) + e^{-i \varphi_1} (00101) + e^{-i \varphi_2} (01000)
+ e^{-i \varphi_3} (10000) \}
\]

(C2)

The second equation shows how, for simplicity, the phase $\varphi_k$ can be realized by a phase $\theta_m$ on the $m$th qubit, as discussed above. Thus each fidelity measurement is parameterized by $\bar{\psi} = \{ \theta_1, \theta_2, \theta_3 \}$, where we have removed $\theta_0$, as it simply gives rise to a global phase. Then, defining $R_{jk} \equiv 2a_j a_k \text{Re}(\rho_{jk}) = \text{Re}(\rho_{jk}) / 2$, each fidelity measurement yields $F_{\bar{\psi}}(\bar{\psi}) = P + C(\bar{\psi})$, where

\[
C(\bar{\psi}) = \sum_{k=j}^d \sum_{j=0}^{d-1} (\bar{\psi}_j \bar{\psi}_k) R_{jk}.
\]

(C3)

Thus to extract $(W_s)$ we must first identify all $R_{jk}$.

For convenience, let us define the sum and difference of two fidelity measurements as $y_{\pm}(\bar{\psi}, \bar{\psi}') = F_{\bar{\psi}}(\bar{\psi}) \pm F_{\bar{\psi}'}(\bar{\psi}')$. Then choosing the inputs of $\bar{\psi} = \{ 001, 010, 100 \}$ and $\bar{\psi}' = \bar{\psi} \oplus \bar{\psi}$ (i.e. the spin-flipped state), we identify three of the six real parts $\{ R_{01}, R_{12}, R_{23} \}$ via the differences of fidelities:

\[
y_{-}(001, 110) = 2(R_{01} + R_{02} - R_{03}),
y_{-}(010, 101) = 2(R_{01} - R_{02} + R_{03}),
y_{-}(100, 011) = 2(-R_{01} + R_{02} + R_{03}).
\]

(C4)

Then, by the sum of fidelities we almost identify the remaining three real parts $\{ R_{12}, R_{13}, R_{23} \}$:

\[
y_{+}(001, 110) = 2(P + R_{12} - R_{13} - R_{23}),
y_{+}(010, 101) = 2(P - R_{12} + R_{13} - R_{23}),
y_{+}(100, 011) = 2(-R_{12} - R_{13} + R_{23}).
\]

(C5)

More specifically, we see that (due to unknown $P$) we need one more measurement, e.g., at $\bar{\psi} = \{ 0, 0, 0 \}$, to solve for the remaining unknowns.

Therefore, a total of $d(d - 1)/2 + 1 = 6 + 1 = 7$ measurements identifies all six real parts of $\rho_{jk}$ and $P$. In
the same manner, \( d(d-1)/2 = 6 \) more measurements at
\( \varphi_k = \{ \pm \frac{\pi}{2} \} \equiv \{ 0, 1 \} \) will identify the imaginary parts of \( \rho_{jk} \),
such that one can reconstruct \( \langle W_j \rangle = \alpha - \max_B [F'(\vec{\psi})] = \alpha - (P + 2 \sum_{j=0}^{d-1} \sum_{k>j} a_j a_k |\rho_{jk}|) \).

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