Engineering mixed states with weak measurements

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It is known that protocols based on weak measurements can be used to steer quantum systems into pre-designated pure states. Here we show that weak-measurement-based steering protocols can be harnessed for on-demand engineering of mixed states. In particular, through a continuous variation of the protocol parameters, one can guide a classical target state to a discorded one, and further on, towards an entangled target state.

Introduction.—A generalized quantum measurement comprises a two-step protocol: (i) switching on, and later off, an interaction Hamiltonian, coupling the quantum system and the quantum detector, leading to a unitary evolution of the combined setup over a prescribed time interval; (ii) performing a projective measurement of the decoupled detector, which leads to a probabilistic quantum jump [1–3]. The detector readout provides information about the system’s state. Measurements are designated strong (projective) or weak, based on the system-detector interaction strength. The former collapses the system to one of the eigenstates of the measured observable. By contrast, generalized (a.k.a. weak) measurements may result in a slight nudge to the system state [4–8]. No matter how weak the measurement is, it always creates an unavoidable impact on the system state through its back-action [9, 10]. Traditionally, this measurement-induced back-action was considered an undesirale effect since the primary purpose of a measurement is to extract information about the system without perturbing it.

Following a disparate paradigm, one may employ the measurement-induced back-action on the system’s state as a means to control the system’s evolution, steering it towards a pre-designated target state [11–15]. A recent work (cf. [16] and references therein) analyzed a host of protocols for engineering pure target states. These, depending on the protocol parameters, may be: (i) non-discorded and non-entangled (“classical”), (ii) discorded and non-entangled, or (iii) discorded and entangled, thus providing us with a smooth navigation tool from classical to fully quantum states.

The guiding principle of our protocol is as follows. Under sufficiently weak measurement, the system evolution can be described by a Lindbladian master equation. Such an equation will have at least one (and in the present case exactly one) steady state that is approached exponentially quickly. We design the protocol such that this steady state is the target state. We first recall a measurement-based protocol whose steady state is an arbitrarily chosen pure state [16]. In order to generate a mixed state, we diagonalize the density matrix of the target state and juxtapose the protocols for stabilizing each of the density matrix’s pure eigenstates. In the explicit example we consider, we find that the rate of converging towards the target state does not significantly depend on its discord or degree of entanglement.

General evolution under repeated blind measurements.—Consider a quantum system in state represented through their discord [18–21]. Discorded quantum states have been proposed as resources for achieving quantum speedup [22, 23], for remote state preparation [24], and for quantum purification protocols [25, 26]. It is thus natural to ask whether a measurement-based steering approach can be used to generate mixed states by design.

Here we introduce a measurement-based protocol which can be used to steer a two-qubit system to an arbitrary predesignated state (pure or mixed), independently of the system’s initial state (the latter is assumed unknown). We illustrate our protocol by considering a family of target states. These, depending on the protocol parameters, may be: (i) non-discorded and non-entangled (“classical”), (ii) discorded and non-entangled, or (iii) discorded and entangled, thus providing us with a smooth navigation tool from classical to fully quantum states.

In order to perform a measurement, the system interacts
with the detector via an interaction Hamiltonian $H_{s-d}$; the joint system-detector state evolves as

$$\rho(t + \tau) = U\rho(t)U^\dagger, \quad (2)$$

where $U = \exp(-iH_{s-d}\tau)$, and $\tau$ is the interaction time. Subsequently the detector state is measured projectively, disentangling the composite system-detector state and generating a measurement back-action on the system state. When discarding (i.e., tracing out) the measurement readouts, a procedure we denote “blind measurement”, the effect of the back-action is represented through

$$\rho_s(t + \tau) = \text{Tr}_d[\rho(t + \tau)]. \quad (3)$$

Following each measurement step, the detector is reset to its initial state $\rho_s^{(0)}$, and then the same measurement procedure is repeated. This protocol gives rise to a non-trivial evolution of the system state. Denoting the system-detector interaction time $\tau$, and taking the continuous time limit $\tau \to dt$, one arrives at the following differential equation for the system state evolution under blind measurements

$$\frac{d\rho_s}{dt} = \mathcal{L}[\rho_s] = i\text{Tr}_d ([\rho(t), H_{s-d}]) - \frac{1}{2}\text{Tr}_d ([H_{s-d}, [H_{s-d}, \rho(t)]]\tau). \quad (4)$$

Here $\mathcal{L}$ is the Liouvillian superoperator acting on the system state, and we dropped the terms $O(\tau^2)$ on the r.h.s. The first term on the right-hand-side of the above equation generates unitary evolution of $\rho_s$, while the second term represents dissipative evolution and can be cast in the form of a Lindbladian. In other words, the above equation can be written as

$$\frac{d\rho_s}{dt} = \mathcal{L}[\rho_s] = i[\rho_s(t), H_s] - \frac{1}{2}\sum_j \left\{ (L_j^\dagger L_j, \rho_s(t)) - 2L_j\rho_s(t)L_j^\dagger \right\}, \quad (5)$$

where $\{,\}$ represents the anti-commutator, $H_s$ is the effective system Hamiltonian, and $L_j$ are the Lindblad jump operators acting on the system state. This way, the sequence of measurements influences the system state in a quasicontinuous manner and ultimately steers it to a steady state determined by the condition

$$\frac{d\rho_s^{(T)}}{dt} = \mathcal{L}[\rho_s^{(T)}] = 0. \quad (6)$$

Steering towards a pure target state.—We now recall the principles of measurement-based steering protocol of Ref. [16], focusing on the two-qubit case, at the center of our analysis. The protocol facilitates stabilizing the system in an arbitrary pure target state $|B_1\rangle$, corresponding to $\rho_s^{(T)} = |B_1\rangle\langle B_1|$.

To implement the protocol we first select three arbitrary states $|B_2\rangle, |B_3\rangle, |B_4\rangle$, and $|B_1\rangle$, such that together with $|B_1\rangle$ they form an orthonormal basis in the four-dimensional Hilbert space of the two-qubit system. Three steps now follow: in the $k^{th}$ step, $k = 1, 2, 3$, a measurement is performed with a system-detector coupling that is designed such that the measurement back-action steers the system away from $|B_{k+1}\rangle$. This is accomplished by choosing $H_{s-d}^k = J (|B_1\rangle\langle B_{k+1}| \otimes \sigma^- + \text{h.c.})$, where the detector is a single qubit acted upon by the Pauli matrices $\sigma^\pm = (\sigma_x^{(d)} \pm i\sigma_y^{(d)})/2$, and $J$ is the coupling strength (for simplicity it is the same in all the three steps). Before each measurement, the detector is initialized in $\rho_d^{(0)} = |\uparrow\rangle\langle\uparrow|$. With the duration of each step being $\tau$, the density matrix evolution over $dt = 3\tau$ is given by

$$\frac{d\rho_s}{dt} = -\frac{g}{2} \sum_{j=1}^4 \left( \{L_j^\dagger L_j, \rho_s\} - 2L_j\rho_s L_j^\dagger \right), \quad (7)$$

where $g = J^2\tau/3$, and the jump operators $L_j = (1 - \delta_{j1})|B_j\rangle\langle B_j|$ (we introduced $L_1 = 0$ for completeness). It follows from Eq. (7) that

$$\frac{d\rho_s}{dt} = 0 \Leftrightarrow \rho_s = \rho_s^{(T)} = |B_1\rangle\langle B_1|, \quad (8)$$

i.e., the system is steered towards the desired pure state.

It is instructive to understand the mechanism by which the above steering works. In each step, the system-detector interaction Hamiltonian $H_{s-d}^k$ steers the system in a two-dimensional subspace spanned by $|B_1\rangle$ and $|B_{k+1}\rangle$ from the state $|B_{k+1}\rangle$ to $|B_1\rangle$ without affecting the rest of $|B_i\rangle$ ($i \neq k + 1$). Since $H_{s-d}^k$ commutes with $|B_1\rangle\langle B_1| \otimes \rho_d^{(0)}$, the measurement does not disturb the system if it is in state $|B_1\rangle$, and the detector then remains in its initial state, $|\uparrow\rangle$. This makes $|B_1\rangle$ not just a steady state of the evolution, but a dark state in the terminology of Ref. [27]: once the system is in $|B_1\rangle$ it is not affected by the detectors. If the system is in $|B_{k+1}\rangle$, a transition to $|B_1\rangle$ (accompanied by the detector state flipping to $|\downarrow\rangle$) happens with probability $\sin^2 J\tau$. Likewise, with probability $\cos^2 J\tau$, the detector does not flip the state and the system remains in $|B_{k+1}\rangle$. Note that if the system is initially in a (coherent or incoherent) superposition of $|B_1\rangle$ and $|B_{k+1}\rangle$, both detector readouts affect the system state: $|\downarrow\rangle$ state of the detector implies the system has jumped from $|B_{k+1}\rangle$ to $|B_1\rangle$, while $|\uparrow\rangle$
implies a change of the weights of the superposition due to different probabilities of the ↑ readout depending on the system state (this is referred in the literature as a “null weak measurement” [28–31] or by a number of different names [32–34]). Averaging over the possible detector readouts and taking the limit \( J\tau \ll 1 \), one obtains the master equation (7).

Steering towards a mixed target state.—We now focus on the key result of the paper: steering the two-qubit system to a desired mixed state. Any mixed target state \( \rho_s^{(T)} \) has a spectral decomposition [35]

\[
\rho_s^{(T)} = \sum_{i=1}^{4} p_i |B_i\rangle\langle B_i|
\]

where \(|B_i=1,...,4\rangle\) form an orthonormal basis in the two-qubit Hilbert space, and \( p_i \geq 0 \) is the probability of the system being in the corresponding \(|B_i\rangle\) state, so that \( \sum_{i=1}^{4} p_i = 1 \). The protocol described in the previous section can be used to steer the system to \(|B_1\rangle\). Furthermore, by exchanging the roles of \(|B_1\rangle\) and one of \(|B_{i=\neq 1}\rangle\), the protocol steers the system to the corresponding \(|B_i\rangle\).

We now show that combining the four protocols, each steering the system to one of \(|B_i\rangle\), with appropriate coupling strengths, \( J \to J_i \), allows to stabilize the mixed state in Eq. (9).

A schematic experimental setup for this complex protocol is presented in Fig. 1. Each part of the protocol, steering the system towards one of the \(|B_i\rangle\) lasts \( 3\tau \). Consequently, the density matrix evolution for \( dt = 4 \times 3\tau = 12\tau \) is described by

\[
\frac{d\rho_s}{dt} = \mathcal{L}[\rho_s] = \sum_{i=1}^{4} \mathcal{L}_i[\rho_s]
\]

\[
= -\frac{1}{2} \sum_{i=1}^{4} g_i \sum_{j=1}^{4} \left( L_j^{(i)} L_j^{(i)^\dagger} \rho_s - 2L_j^{(i)^\dagger} L_j^{(i)} \rho_s L_j^{(i)^\dagger} \right)
\]

where \( g_i = J_i^2 \tau /12 \) and \( L_j^{(i)} = (1-\delta_{ij})|B_i\rangle\langle B_j| \). Equation (10) has a unique steady state,

\[
\rho_s = \rho_s^{(T)} = \frac{1}{\sum_{i=1}^{4} g_i} \left( \sum_{j=1}^{4} g_j |B_j\rangle\langle B_j| \right).
\]

We note that choosing

\[
g_j = \bar{g} p_j ,
\]

where \( j = 1,2,3,4 \) and \( \bar{g} = \sum_{i=1}^{4} g_i \), stabilizes the desired target state in Eq. (9). Therefore, measurement-based steering towards an arbitrary mixed state requires diagonalizing the density matrix of the latter, bringing the state to the form (9), and concurrent utilization of the four protocols, each stabilizing one of the pure eigenstates which make up the mixed-state density matrix.

We emphasize that the simple-looking result in Eq. (12) is highly non-trivial. Indeed, for arbitrary four Lindbladians \( \mathcal{L}_i \) such that \( \mathcal{L}_i[|B_j\rangle\langle B_j|] = 0 \), one cannot guarantee that \( \sum_i g_i \mathcal{L}_i[\sum_j g_j |B_j\rangle\langle B_j|] = \sum_{i,j} g_i g_j \mathcal{L}_i[|B_j\rangle]\langle B_j||B_j\rangle\langle B_j|] = 0 \). The fact that it is true in our case is a special feature of the Lindbladians involved in our measurement-based steering protocol.

If we record the detector’s readouts, the dynamics of the system subject to our protocol can be understood as follows: The limit of weak measurement addressed here, \( J\tau \ll 1 \), implies that a detector click (one of the detectors flipping its state to \(|\downarrow\rangle\) during the measurement) is a rare event. It follows that the system, initially prepared in a possibly coherent superposition of different \(|B_i\rangle\), will experience coherence-preserving dynamics induced by a “no-click” back-action for a long time \( \sim \bar{g}^{-1} \) (this is akin to “null weak measurements” [28–31]). Eventually, a click will register, bringing the system to one of the \(|B_i\rangle\) states and destroying coherence between different \(|B_i\rangle\) components. From this point on, the no-click dynamics does not affect the system state, while rare clicks make it jump between different \(|B_i\rangle\) states. The system spending random amounts of time in different \(|B_i\rangle\) results in the average state given by Eq. (9). Notably, the system, while being on average in the target state, probabilistically jumps among the constituent pure states \(|B_i\rangle\), which is manifested by occasional detector clicks. In other words, a mixed target state in our protocol is a steady state but not a dark state.

Steering to “classical” vs “quantum” states.—We now illustrate our protocol with an example where the system can be steered into a family of states which can be classical or quantum depending on the measurement couplings employed. Consider the following two-qubit state,

\[
\rho = p_1 |\uparrow\uparrow\rangle\langle \uparrow\uparrow| + p_2 |\psi^+\rangle\langle \psi^+| + p_3 |\psi^-\rangle\langle \psi^-| + p_4 |\downarrow\downarrow\rangle\langle \downarrow\downarrow|.
\]
and Peres-Horodecki criterion \cite{36, 37}, one shows that the initial state can be pure (\(\rho_i(0) = \frac{1}{4} \sum_{i,j} |B_i\rangle\langle B_j|\), top) or mixed (\(\rho_i(0) = \frac{1}{2} (|B_1\rangle\langle B_1| + |B_2\rangle\langle B_2|)\), bottom). \(\bar{F} = 1 - \left(\text{Tr} \sqrt{\rho(t)\rho(t)}\right)^2\) is the deviation from perfect fidelity of the target state preparation. Irrespective of the initial state, the system has essentially converged to the target state with \(\alpha = \beta = \frac{1}{2}\) at \(\bar{g}t = 5\).

Figure 3. The density matrix of the two-qubit system as it is steered towards a target state, \(\bar{\rho}\), of Eq. (13). The magnitudes of the density matrix entries in the basis \(|B_i\rangle\) (17) are represented by a color scale. The initial state can be pure \((\rho_i(0) = \frac{1}{4} \sum_{i,j} |B_i\rangle\langle B_j|\), top) or mixed \((\rho_i(0) = \frac{1}{2} (|B_1\rangle\langle B_1| + |B_2\rangle\langle B_2|)\), bottom). \(\bar{F} = 1 - \left(\text{Tr} \sqrt{\rho(t)\rho(t)}\right)^2\) is the deviation from perfect fidelity of the target state preparation. Irrespective of the initial state, the system has essentially converged to the target state with \(\alpha = \beta = \frac{1}{2}\) at \(\bar{g}t = 5\).

where \(|\psi\rangle = \frac{1}{\sqrt{2}} (|\uparrow\downarrow\rangle \pm |\downarrow\uparrow\rangle)\), and

\[
  p_1 = \frac{(1 - \beta + \alpha(1 + \beta))}{4}, \quad p_2 = \frac{(1 - \alpha)(1 - \beta)}{4},
  
  p_3 = \frac{(1 - \alpha)(1 + 3\beta)}{4}, \quad p_4 = \frac{(1 - \beta + \alpha(1 + \beta))}{4},
\]

Here \(\alpha\) and \(\beta\) are two independent parameters such that \(0 \leq \alpha, \beta \leq 1\), and the coefficients \(p_i\) correspond to the \(p_i\) in Eq. (9). This state may or may not have quantum correlations depending on \(\alpha\) and \(\beta\). For example, using Peres-Horodecki criterion \cite{36, 37}, one shows that \(\bar{\rho}\) is separable (not entangled) if and only if \(\alpha \geq \frac{3\beta - 1}{3\beta + 1}\).

Quantum correlations are commonly quantified via concurrence (a measure of entanglement) \cite{20, 38} and quantum discord \cite{18–21}. Calculating discord \(\mathcal{D}\) and concurrence \(\mathcal{C}\) for an arbitrary state is a challenging task, but for the state in Eq. (13) both of them can be calculated analytically \cite{20, 38–40}. They are respectively given by

\[
\begin{align*}
\mathcal{Q}(\bar{\rho}) &= \frac{1 - \alpha}{4} \left( (1 - \beta) \log_2[(1 - \alpha)(1 - \beta)] \\
&\quad - 2(1 + \beta) \log_2[(1 - \alpha)(1 + \beta)] \\
&\quad + (1 + 3\beta) \log_2[(1 - \alpha)(1 + 3\beta)] \right),
\end{align*}
\]

and

\[
\mathcal{C}(\bar{\rho}) = \begin{cases} 
\frac{3\beta(1 - \alpha) - (1 + \alpha)}{4} & \text{for } \alpha < \frac{3\beta - 1}{3\beta + 1}, \\
0 & \text{otherwise}.
\end{cases}
\]

Note that the discord \(\mathcal{Q}(\bar{\rho})\) is only zero when: (i) \(\alpha = 1\) or (ii) \(\beta = 0\), while the concurrence \(\mathcal{C}(\bar{\rho})\) exhibits a sharp change of behavior at finite \(\alpha\) and \(\beta\), cf. Fig. 2. The state \(\bar{\rho}\) can thus be purely classical (both discord and concurrence vanish), discorded (discord vanishes), or entangled (discord and concurrence are both non-zero) depending on the parameters \(\alpha\) and \(\beta\).

We may generate \(\bar{\rho}\) using the protocol described above. First, the target state density matrix should be diagonalized. It is evident from Eq. (13) that the eigenbasis of \(\bar{\rho}\) is

\[
|B_1\rangle = |\uparrow\rangle, \quad |B_2\rangle = |\psi^+\rangle, \quad |B_3\rangle = |\psi^-\rangle, \quad |B_4\rangle = |\downarrow\rangle.
\]

Using Eq. (12), one obtains that the couplings \(g_i\) in Eq. (10) are \(g_i = \bar{g} p_i\), where \(\bar{g} = \sum_i g_i\) characterizes the total strength of all measurements employed. Figure 3 illustrates the time evolution of the two-qubit system as it is steered from an initial state to the target state \(\bar{\rho}\). The deviations from the target state decay exponentially in time; the decay rates are determined by the real parts of the non-zero eigenvalues of the Liouvillian superoperator \(\mathcal{L}\), cf. Eq. (10). The smallest (in magnitude) real part determines the slowest convergence rate; its dependence on the target state is presented in Fig. 4. Note that the convergence rate does not depend significantly on \(\alpha\) and \(\beta\), implying that our protocol works equally well for both entangled and non-entangled states, as well as for states in the vicinity of the transition between the two regions.

Discussion.—We have proposed a measurement-based protocol that can generate any two-qubit state by design (pure or mixed), starting from an arbitrary unknown initial state. We illustrate the protocol with an example, in which the target state can be classical, discorded, or entangled, depending on relative strengths of the measurements employed in the protocol.

We emphasize a conceptual difference of our proto-
col to more conventional drive-and-dissipation schemes [27, 41–52], where environment is employed to relax the system to a desired state. The two crucial distinctions here are: (a) the relaxation is induced by measurements, implying the possibility to use the measurement readouts to confirm the system’s desired behavior (and, possibly, hasten the convergence towards the target state); (b) our system does not have a Hamiltonian (no “drive”).

Finally, we note that our protocol is, in principle, generalizable to $N$-qubit systems. In practice this, however, is quite tricky. First, the number of measurements employed in our protocol is proportional to the Hilbert space size. Second, each such measurement would, in general, involve coupling a detector to all $N$ qubits, hence requiring $N + 1$-body interactions. It is, therefore, interesting to develop scalable alternatives to our protocol (possibly with a restricted class of states that can be stabilized, e.g., matrix product operators).

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