The relation between the quantum discord and quantum teleportation: The physical interpretation of the transition point between different quantum discord decay regimes

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Abstract – We study quantum teleportation via Bell-diagonal mixed states of two qubits in the context of the intrinsic properties of the quantum discord. We show that when the quantum-correlated state of the two qubits is used for quantum teleportation, the character of the teleportation efficiency changes substantially depending on the Bell-diagonal-state parameters, which can be seen when the worst-case-scenario or best-case-scenario fidelity is studied. Depending on the parameter range, one of two types of single-qubit states is hardest/easiest to teleport. The transition between these two parameter ranges coincides exactly with the transition between the range of classical correlation decay and quantum correlation decay characteristic for the evolution of the quantum discord. The correspondence provides a physical interpretation for the prominent feature of the decay of the quantum discord.

Introduction. – The quantum discord [1–3], a measure of bi- and multi-partite quantum correlations, has recently attracted much attention. This is due to the fact that the discord indicates the existence of quantum correlations in many partially mixed states which have no entanglement present. Specifically, the quantum discord, $D$, does not display any sudden-death–type phenomenon, since the zero-discord states form a set of measure zero [4]. This subset of states is the set of all truly classical states chosen asymptotically by decoherence [5,6]. Hence, a smooth, continuous decoherence process cannot lead to a sudden and continued disappearance of quantum correlations mid-evolution (before a fully mixed, completely dephased state is reached). This suggests that the sudden death of entanglement signifies not, as was previously believed, the disappearance of all quantum correlations, but the crossing of a threshold of a given, small amount of correlations and the disappearance of quantum correlations of a certain type. Although below this threshold many quantum informational tasks are no longer possible, methods of performing quantum computation on zero-entanglement states for which the quantum discord is non-zero have already been devised [7–11]. Furthermore, it has been very recently shown experimentally how entanglement can be shared between distant parties via non-entangled states with non-zero discord [12–16].

Unfortunately, computing the quantum discord for an arbitrary density matrix is an extremely involved task even in the simplest two-qubit case [17]. This led to the emergence of the geometric quantum discord [18], which is defined as the minimal Hilbert-Schmidt distance of a given state from the set of zero-discord states. Although an explicit formula for the geometric discord given a two-qubit density matrix does not yet exist (one that does not require minimization over the set of all zero-discord states), such formulas exist for the lower [18] and upper [19] bounds on the geometric discord. The geometric discord is a good measure to distinguish between zero-discord and non-zero-discord states, but because of the properties of the Hilbert-Schmidt distance, it is not a good measure for the amount of quantum correlations present in a given state. In fact, because the Hilbert-Schmidt distance is sensitive to the purity of the state, the geometric discord may be increased by non-unitary evolution of a single qubit (the unmeasured one) [20–22], which should not increase inter-qubit quantum correlations. One solution to
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this problem is rescaling the geometric discord following ref. [23]; a lot of work is also being done to find explicit formulas for discord measures utilizing more appropriate distance measures [24–27].

Regardless of these difficulties, there are properties of typical discord evolutions, $D(t)$, under decoherence processes that have already become apparent. Perhaps the most prominent and most baffling is the behavior of the geometric quantum discord and rescaled geometric quantum discord [23], and the transition was shown to occur at exactly the same point as the transition in the decay of $D(t)$ [31,32]. The study of the geometric discord revealed similar transitions occurring for non-Bell-diagonal initial states and evolutions.

The above discussion suggests that the transition between the “classical” and “quantum” decay regimes is a rather general feature of $D$ dynamics. There is, however, no intuitive understanding of the possible operational meaning of this transition, which is not surprising in the light of the fact that the operational significance of the quantum discord is a subject of ongoing research [33,34]. In the following, we give a simple interpretation of an observable (and practically relevant) quantity that changes discontinuously at the transition point of discord dynamics.

We investigate the quantum teleportation protocol with the entangled Bell state used for teleportation being subject to decoherence. It is well known that the time at which the average teleportation fidelity becomes smaller than 2/3 is the time at which the bipartite state used as a resource becomes separable [35–37]. Here we focus on teleportation fidelity minimized over the teleported states (i.e. the worst-case-scenario fidelity) or maximized over the teleported states (i.e. the best-case-scenario fidelity; the two situations display an equivalent transition), and we show that at the time of the transition between the two regimes of quantum discord decay, the nature of the state for which teleportation has the lowest/highest fidelity changes.

Classical and quantum decoherence regimes of quantum discord dynamics. — The quantum discord is defined as the difference between two classically equivalent formulas for mutual information [1]. The formula which is referred to as mutual information in ref. [29] is given by [1]

$$I(\rho_{AB}) = S(\text{Tr}_B \rho_{AB}) + S(\text{Tr}_A \rho_{AB}) - S(\rho_{AB}),$$

where the von Neumann entropy is given by $S(\rho) = -\text{Tr} \rho \log_2 \rho$. This quantity was generalized from the classical language of probability distributions in a straightforward manner to the language of density matrices, while the Shannon entropy was replaced by the von Neumann entropy. The other formula for classical mutual information, which is in the quantum context referred to as classical correlations (in ref. [29]), cannot be generalized in a direct manner, because the classical formula involves conditional entropy,

$$C(A : B) = H(A) - H(A|B),$$

where $H$ denotes the Shannon entropy and $A$ and $B$ are random variables. Conditional entropy $H(A|B)$ requires the specification of the state of $A$ given the state of $B$, which in quantum mechanics is ambiguous until the measurement performed on $B$ is specified. Hence, the conditional von Neumann entropy can be found given the complete measurement on subsystem $B$ and the resulting formula for classical correlations is [1,2]

$$C(\rho_{AB}) = \max_{\{\Pi_k\}} \left[ S(\text{Tr}_B \rho_{AB}) - S(\rho_{AB}|[\Pi_k]) \right],$$

where $\{\Pi_k\}$ is a complete set of orthonormal projective operators corresponding to a von Neumann measurement of subsystem $B$. The index $k$ denotes the outcome of a given measurement and the formula involves maximization over the set of projective measurements. Therefore, the formula of eq. (2) yields the information gained about the system $A$ after the measurement $\{\Pi_k\}$ on system $B$. The quantum discord of a given state $\rho_{AB}$ is then given by

$$D(\rho_{AB}) = I(\rho_{AB}) - C(\rho_{AB}).$$

Following ref. [29], the value of the quantum discord can be found for any Bell-diagonal two-qubit state. If the Bell-diagonal two-qubit density matrix is written in the form

$$\rho_{AB} = \begin{pmatrix} \rho_{\Phi} & 0 & 0 & \sigma_{\Phi} \\ 0 & \rho_{\Phi} & \sigma_{\Phi} & 0 \\ 0 & \sigma_{\Phi} & \rho_{\Phi} & 0 \\ \sigma_{\Phi} & 0 & 0 & \rho_{\Phi} \end{pmatrix},$$

where all four parameters are real, the mutual information is given by

$$I(\rho_{AB}) = 2 + \sum_{i=\Psi,\Phi} [(\rho_i + \sigma_i) \log_2(\rho_i + \sigma_i)$$

$$+ (\rho_i - \sigma_i) \log_2(\rho_i - \sigma_i)].$$

Note that any continuous and differentiable evolution of $\rho_{AB}$ that retains the Bell-diagonal form of eq. (4)
must lead to a continuous and differentiable evolution of $I(\rho_{AB})$. An explicit formula for the classical correlations of eq. (2) can also be found for any Bell-diagonal state [28]. These correlations are given by

$$C(\rho_{AB}) = \frac{1}{2}[(1 + \chi) \log_2(1 + \chi) + (1 - \chi) \log_2(1 - \chi)],$$

where $\chi = \max\{|\Delta|, |\sigma_\Phi| + |\sigma_\bar{\Phi}|\}$, in which we used $\Delta = \rho_{\Phi} - \rho_{\bar{\Phi}}$. The maximization allows for indistinguishability points in the evolution of classical correlations which occurs when the plane

$$|\Delta| = |\sigma_\Phi| + |\sigma_\bar{\Phi}|$$

is transgressed, resulting in indistinguishability points of the quantum discord which is given by the difference of the smooth mutual information and the indistinguishable classical correlation function, eq. (3). Note that under a pure dephasing decoherence process, the classical correlation function remains constant in the $|\Delta| > |\sigma_\Phi| + |\sigma_\bar{\Phi}|$ regime, which we will call the quantum decoherence regime following ref. [29], while it decays in the classical decoherence regime, $|\Delta| < |\sigma_\Phi| + |\sigma_\bar{\Phi}|$. Below we consider decoherence which leads to $C(t)$ and $D(t)$ both decaying in the two regimes defined by the above inequalities.

**Relation between teleportation fidelity and discord.** – To gain some understanding of the physical significance of this transition, let us turn to the quantum teleportation of an unknown qubit state by means of an entangled two-qubit state [38] in the situation when the entangled state has previously undergone partial decoherence. We will focus on the scenario where the initial entangled state is the $|\Phi^+\rangle = 1/\sqrt{2}(|00\rangle + |11\rangle)$ Bell state, corresponding to $\rho_{\Phi} = \sigma_\Phi = 1/2$ and $\rho_{\bar{\Phi}} = \sigma_{\bar{\Phi}} = 0$ in eq. (4); the assumption is made for simplicity; the same results are acquired regardless of the chosen Bell state. For such an initial state, decoherence often retains the Bell-diagonal form, but rarely provides the means to induce coherences between the $|01\rangle$ and $|10\rangle$ components ($\sigma_\Phi \neq 0$). It can, however, induce non-zero occupations corresponding to $|01\rangle$ and $|10\rangle$ [39,40], and possibly change the diagonal matrix elements in such a way that the state ceases to be Bell-diagonal. Since we want to maintain the Bell-diagonal form at all times (so that we can analytically calculate the quantum discord), below we will focus on a physically well-motivated example of decohering channel which preserves this form with $\sigma_\Phi = 0$.

First let us explain the connection between the two regimes of discord dynamics and the teleportation fidelity. The unknown qubit to be teleported is $|\psi\rangle = \alpha|0\rangle + \beta|1\rangle$, with $|\alpha|^2 + |\beta|^2 = 1$. After teleporting $|\psi\rangle$ using the state of eq. (4) with $\sigma_\Phi = 0$ instead of the $|\Phi^+\rangle$, the teleported state is equal to

$$\rho_t = |0\rangle\langle 0| (2\rho_\Phi|\alpha|^2 + 2\rho_\Phi|\beta|^2) + |1\rangle\langle 1| (2\rho_{\bar{\Phi}}|\alpha|^2 + 2\rho_{\bar{\Phi}}|\beta|^2) + |0\rangle\langle 2\sigma_\Phi \alpha \beta^* + \text{H.c.}$$

and the fidelity of teleportation is given by

$$F^2 = \langle \psi | \rho_t | \psi \rangle = 2\rho_{\bar{\Phi}} - 4(\Delta - |\sigma_{\Phi}|)(1 - |\alpha|^2).$$

The standard procedure now is to find the average fidelity of the teleported state (averaged over all possible states to be teleported) and compare it to the classical limit of teleportation capability which is equal to 2/3 [35–37]. The average fidelity is given by

$$F^2_{av} = 2\rho_{\bar{\Phi}} - \frac{2}{3}(\Delta - |\sigma_{\Phi}|),$$

which is a smooth function of time. In order to see the correspondence between the transition in the evolution of quantum correlations and teleportation capability, one should rather look at the extremal scenarios, namely, the fidelity of the teleported state minimized or maximized over all the $|\psi\rangle$ states. Analyzing the extrema of eq. (9) as a function of $|\alpha|^2 \in [0,1]$ we find that for $|\Delta| < |\sigma_{\Phi}|$ (in the classical decoherence regime), the minimal fidelity $F^2_{min} = 1/2 + \Delta$ occurs for $|\psi\rangle = |0\rangle, |1\rangle$ (the poles on the Bloch sphere), and the maximal fidelity $F^2_{max} = 1/2 + \sigma_{\Phi}$ is obtained for $|\psi\rangle = 1/\sqrt{2}(|0\rangle + \exp(i\phi)|1\rangle)$ (the states on the equator of the Bloch sphere). On the other hand, for $|\Delta| > |\sigma_{\Phi}|$ (in the quantum decoherence regime), the states which are easiest/hardest to teleport (with the corresponding fidelities) trade places, so that we always have $F^2_{min} = 1/2 + \min(\Delta, \sigma_{\Phi})$ and $F^2_{max} = 1/2 + \max(\Delta, \sigma_{\Phi})$. At the transition point, $\Delta = \sigma_{\Phi}$, the fidelity from eq. (9) is independent of the state $|\psi\rangle$, so that it is equal to its average value, $2\rho_{\bar{\Phi}} = \Delta + 1/2$. If $\Delta > 1/6$ at this point in time, the transition between the two regimes of quantum discord dynamics occurs when the state is still entangled, and $F^2_{av}$ is larger than its maximal classical value.

The character of state $|\phi\rangle$ which is easiest/hardest to teleport in a given regime can be also connected with the character of the classical states closest to $\rho_{AB}$ [41]. In the “classical” regime there are two such states, $\rho_\alpha$ and $\rho_\beta$ given by

$$\rho_\alpha = \frac{1}{4} + \frac{\sigma_\Phi}{2} \langle aa | aa \rangle + \langle \bar{a} \bar{a} | \bar{a} \bar{a} \rangle,$$

$$\rho_\beta = \frac{1}{4} - \frac{\sigma_\Phi}{2} \langle aa | aa \rangle + \langle \bar{a} \bar{a} | \bar{a} \bar{a} \rangle,$$

where $a = x, y$, and the states $|a\rangle$ and $|\bar{a}\rangle$ are eigenstates of the Pauli matrices $\sigma_\alpha$ with $\pm 1$ eigenvalues. Such a classically correlated state $\rho_a$ allows for “teleportation” of $|\phi\rangle = |a\rangle, |\bar{a}\rangle$ with fidelity of $1/2 + \sigma_\Phi$. The fact that these are the closest classical states in the $\sigma_{\Phi} > \Delta$ regime gives an intuitive explanation for the maximal fidelity being achieved for the $|\phi\rangle$ state having its Bloch vector in the $xy$ plane. On the other hand, in the “quantum” regime the closest classical state is given by a formula analogous to the one from eq. (11), but with $a = z$ and $\sigma_\Phi$ being replaced by $\Delta$. The correlations in this state allow for classical enhancement of teleportation fidelity of $|0\rangle$ and $|1\rangle$ states.
The zero-discord line is the boundary that separates classical decoherence from entangled decoherence. The red dashed line is the indifferentiability plane given by eq. (7), which is a line where the quantum system cannot be distinguished from a classical one. The blue line denotes the classical decoherence regime (yellow triangle). The uncolored region corresponds to parameter values that yield an unphysical density matrix. The red dashed line is the zero-discord line (entanglement vanishes below the unphysical density matrix). The dotted lines are isoliscors, corresponding to $D = 0.1, 0.3, 0.5, 0.7$ going from bottom to top.

Figure 1 shows the regimes of decoherence for Bell-diagonal density matrices with $\sigma_\Phi = 0$ with respect to the amplitude of the coherence present in the system and the difference of the two distinct two-qubit occupations. The border between the region of quantum decoherence (cyan) and classical decoherence (yellow) in terms of the quantum discord is denoted by the blue line. The Bell state used above is located in the upper right corner of the figure ($\Delta = 1/2$ and $|\sigma_\Phi| = 1/2$) and its decoherence will move it down and/or to the left. In fact, a pure dephasing process of an initial Bell state results in its moving straight down along the $\Delta = 1/2$ line, which is located solely in the quantum decoherence regime. Vanishing entanglement corresponds to the point when $F^2_{\sigma_\Phi}$ reaches the classical communication limit of $2/3$ [35–37]. This boundary is denoted by a red, dashed line in fig. 1.

As seen, sudden death of entanglement [42–44] (i.e. crossing of the separability boundary while $\sigma_\Phi$ is nonzero) can be achieved both via the quantum decoherence regime and the classical decoherence regime. Moreover, entanglement sudden death can occur for higher values of coherence $\sigma_\Phi$ in the “classical” parameter range, because it is susceptible to the disturbance of qubit occupations [45]. Incidentally, the zero-discord line is the $|\sigma_\Phi| = 0$ axis and coincides with the minimized teleportation fidelity reaching the minimal value of $1/2$.

It is straightforward to generalize the teleportation procedure via a decohered two-qubit state to any Bell-diagonal state, eq. (4), and to show that the discord transition plane of eq. (7) still corresponds to a transition in the worst-case-scenario teleportation from the situation when the states on the pole of the Bloch sphere are hardest to teleport, and when states on the equator of the Bloch sphere are hardest to teleport. Hence the relation between quantum teleportation and the quantum discord presented in this section holds for any state of the form given by eq. (4).

**Possible physical realization.** Let us now present a model of open-system dynamics in which the above-discussed transition between two teleportation regimes occurs at finite time $t_\omega$. It is clear that the classical decoherence requires a minimal amount of dephasing (decay of $\sigma_\Phi$) consistent with a given amount of relaxation (decay of $\Delta$), suggesting the use of an amplitude damping channel. In order for the state to remain Bell-diagonal, we use the generalized amplitude damping channel [39] corresponding to an environment being in equilibrium at temperature much higher than the energy splitting of the qubits. Assuming Markovian decoherence, we have $\Delta(t) = \Delta(0)e^{-\Gamma t}$ and $\sigma_\Phi(t) = \sigma_\Phi(0)e^{-\Gamma t/2}$, where $\Gamma = \Gamma_A + \Gamma_B$, with $\Gamma_A, \Gamma_B$ being the longitudinal relaxation rates of the two qubits forming the entangled state used for teleportation.

However, under the action of this channel alone, the evolution starting form one of the Bell states remains in the classical regime at all times. In order for the transition to the quantum decoherence regime to occur, the additional dephasing process, having negligible contribution at short times, has to become stronger at longer times. This will happen when the qubits forming the entangled state used for teleportation are coupled via their $\sigma_z$ operators to a slowly fluctuating bath. Low-frequency ($\omega \ll k_BT$), where $T$ is the bath temperature) environmental fluctuations can be treated as classical noise $\xi(t)$ [46], which, under a realistic assumption of Gaussian statistics, is described by a spectral density $S(\omega) \equiv \int e^{i\omega t}\langle \xi(t)\xi(0)\rangle d\omega$ (where $\langle \ldots \rangle$ denotes averaging over noise realizations).

When most of spectral weight is concentrated at low frequencies ($\omega \ll 1/T_2$, where $T_2$ is the characteristic dephasing timescale), the total noise power is $\sigma^2 \propto F^2_{\sigma_\Phi} S(\omega) d\omega$. This is a typical situation in many solid-state–based qubits (with $1/f$ charge and flux noise [47] and hyperfine interaction with nuclear bath [40,48,49] being prominent examples). Such a noise leads to an additional decay of $\sigma_\Phi(t)$, given by $\sigma_\Phi(t) \rightarrow \sigma_\Phi(t) \exp[-(\gamma t)^2]$, with $\gamma = \sigma/\sqrt{2}$. Note that since the dephasing and relaxation are caused by very different processes (the former by low-energy fluctuations in the environment, the latter by high-energy processes involving exchanges of energy quanta corresponding to qubits’ energy splittings, which are typically much larger than $k_BT$), treating them as independent and additive processes is a reasonable approximation. Decay of quantum correlations due to such low-frequency or quasi-static baths was a subject of a few recent works [32,50–54]. Within this model of decoherence we obtain a transition between two regimes of quantum discord decay at $t = \Gamma/2\gamma^2$, at which point we have $F^2_{\sigma_\Phi} = \frac{1}{2}(1 + \exp[-(\frac{1}{2}(\Gamma/\gamma)^2])$, so that this worst-case scenario fidelity still exhibits entanglement-related enhancement when $\Gamma/\gamma < \sqrt{2/\ln 3}$. The modulus of the
discontinuity of the derivative of $D(t)$ at $t = t_c$ is given by 
\[ \frac{1}{2} \Gamma \Delta_c \ln \left( \frac{1 + \Delta^2_c}{1 - \Delta^2_c} \right), \]
where $\Delta_c = \frac{1}{2} \exp\left(-\Gamma/\gamma\right)$. The quantum-classical transition is thus most visible for low values of $\Gamma/\gamma$, which means that the rms $\sigma$ of the low-frequency phase noise should be much larger than the energy relaxation rate of the qubits. When the energy relaxation is in fact caused by transverse coupling to high-frequency noise, the power of this noise at frequencies corresponding to qubits’ energy splittings (which is proportional to $\Gamma$) should be much smaller than $\sigma$. When the low and high-frequency noises have the same physical origin, this requirement is consistent with our assumption of the dominantly low-frequency character of the environmental fluctuations.

Note that while the above result has been obtained in the case of pure dephasing coupling to low-frequency noise (which leads to Gaussian decay of $\sigma \phi(t)$), a qualitatively analogous transition between classical and quantum decoherence regimes can occur for more general coupling to low-frequency noise. For example, for transverse coupling to noise in the presence of large energy splitting (i.e., the qubit Hamiltonian $H = \Omega \sigma_z/2 + \xi(t) \sigma_x/2$ with $\Omega \gg \sqrt{\xi^2(t)}$), the evolution of the qubits is approximately [55–59] of the pure dephasing form (only with the term quadratic in the noise, $\xi^2(t)/2\Omega$, contributing to the qubits’ splittings), and the initial coherence decay is $\sigma \phi(t) \approx 1 - (t/\gamma)^2$, with $\gamma' \approx \gamma^2/\Omega$. At very short times $(t/\gamma')^2 \ll \Gamma t$, the discord dynamics will be in the classical regime, while at longer times the quantum decoherence regime can be entered. Interestingly, since the asymptotic decay of $\sigma \phi(t)$ in this case is a power law [50,55,59] (so that for long $t$ again we have $\sigma \phi(t) > \Delta(t)$), the transition into the quantum decoherence regime has to be followed by a re-entry into the classical decoherence regime at a later time.

**Conclusion.** – We have studied the properties of quantum correlations present in mixed two-qubit states of Bell-diagonal form. We have shown that there is a direct correspondence between the transition point between the regime of classical correlation decay and the regime of quantum correlation decay, which is a characteristic of the evolution of the quantum discord, and the transition point which appears for the worst-case-scenario fidelity of quantum teleportation. For teleportation, the two regions being transgressed correspond to two different classes of states which are hardest to teleport. Those are either the equal superposition states in the quantum decoherence regime, or the single-component $|0\rangle/|1\rangle$ states in the classical decoherence regime. This shows that there is a qualitative physical difference between the quantumly correlated states studied, depending on the correlation decay regime in which the state is located. Furthermore, the fact that the transition point seen is also by studying the geometric discord and that similar transition points are seen for initial states which for technical reasons cannot be handled using the regular quantum discord, suggests that there is an underlying physical meaning to all such transition points, and that classical-quantum decoherence transition is a widely occurring property of the quantum discord.

**REFERENCES**

[1] Olliver H. and Žurek W., Phys. Rev. Lett., 88 (2001) 017901.
[2] Henderson L. and Vedral V., J. Phys. A: Math. Theor., 34 (2001) 6899.
[3] Modi K., Brodutch A., Cable H., Paterek T. and Vedral V., Rev. Mod. Phys., 84 (2012) 1655.
[4] Ferraro A., Aolita L., Cavalcanti D., Cucchiettti F. M. and Acín A., Phys. Rev. A, 81 (2010) 052318.
[5] Žurek W. H., Rev. Mod. Phys., 75 (2003) 715.
[6] Horneberger K., Lect. Notes Phys., 768 (2009) 221.
[7] Meurer B., Heitmann D. and Flego K., Phys. Rev. Lett., 68 (1992) 1371.
[8] Knill E. and Laflamme R., Phys. Rev. Lett., 81 (1998) 5672.
[9] Dat A., Shi A. and Caves C., Phys. Rev. Lett., 100 (2008) 050502.
[10] Passante G., Moussa O., Trotter D. A. and Laflamme R., Phys. Rev. A, 84 (2011) 044302.
[11] Horodecki P., Tuziemski J., Mazurek P. and Horodecki R., Phys. Rev. Lett., 112 (2014) 140507.
[12] Dakić B., Lipp Y. O., Ma X., Ringbauer M., Kropatschek S., Barz S., Paterek T., Vedral V., Zeilinger A., Brukner C. and Walther P., Nat. Phys., 8 (2012) 666.
[13] Silberhorn C., Phys. Rev., 6 (2013) 132.
[14] Fedrizzi A., Zuppardo M., Gillett G. G., Broome M. A., Almeida M. P., Paternostro M., White A. G. and Paterek T., Phys. Rev. Lett., 111 (2013) 230504.
[15] Peuntinger C., Chille V., Mišta L., Korolková N., Förltsch M., Korger J., Marquardt C. and Leuchs G., Phys. Rev. Lett., 111 (2013) 230506.
[16] Vollmer C. E., Schulze D., Eberle T., Händchen V., Fiurášek J. and Schindl R., Phys. Rev. Lett., 111 (2013) 230505.
[17] Hu Y., New J. Phys., 16 (2014) 033027.
[18] Dakić B., Vedral V. and Brukner Č., Phys. Rev. Lett., 105 (2010) 190502.
[19] Miranowicz A., Horodecki P., Chhajlany R. W., Tuziemski J. and Sperling J., Phys. Rev. A, 86 (2012) 042123.
[20] Piani M., Phys. Rev. A, 86 (2012) 034101.
[21] Tufarelli T., Girolami D., Vasile R., Bose S. and Adesso G., Phys. Rev. A, 86 (2012) 052326.
[22] Hu X., Fan H., Zhou D. L. and Liu W.-M., Phys. Rev. A, 87 (2013) 032340.
[23] Tufarelli T., MacLean T., Girolami D., Vasile R. and Adesso G., J. Phys. A, 46 (2013) 275308.
