Quantum Measurement Bounds beyond the Uncertainty Relations

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In quantum mechanics, the Heisenberg uncertainty relations and the Cramér-Rao inequalities typically limit the precision in the estimation of a parameter through the standard deviation of a conjugate observable. Here we extend these relations by giving a bound to the precision of a parameter in terms of the expectation value of the conjugate observable. This has both foundational and practical consequences: in quantum optics it resolves a controversy over which is the ultimate precision in interferometry.

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Quantum mechanics limits the accuracy with which one can measure conjugate quantities: the Heisenberg uncertainty relations [1,2] and the quantum Cramér-Rao inequality [3–6] show that no procedure for estimating the value of some quantity (e.g., a relative phase) can have a precision that scales more accurately than the inverse of the standard deviation of a conjugate quantity (e.g., the energy) evaluated on the state of the probing system. This paper exhibits a new bound on quantum measurement: we prove that the precision of measuring a quantity cannot scale better than the inverse of the total number of photons employed in the interferometer.

The statistical nature of quantum mechanics induces fluctuations that limit the ultimate precision that can be achieved in the estimation of any parameter $x$. These fluctuations can be connected to the properties of a conjugate operator $H$ that generates translations $U_x = e^{ixH}$ of the parameter $x$. In particular, if the encoding stage is repeated several times using $\nu$ identical copies of the same probe input state $\rho_x$, the root mean square error (RMSE) $\Delta X$ of the resulting estimation process is limited by the quantum Cramér-Rao bound [3–6]

$$\Delta X \geq 1/\sqrt{\nu Q(x)}$$

where $Q(x)$ is the quantum Fisher information. For pure probe states and unitary encoding mechanism $U_x$, $Q(x)$ is equal to the variance $(\Delta H)^2$ (calculated on the probe state) of the generator $H$. In this case, the Cramér-Rao bound takes the form

$$\Delta X \geq 1/(2\sqrt{\nu \Delta H}) \quad (1)$$

of an uncertainty relation [5,6]. This bound is asymptotically achievable in the limit of $\nu \rightarrow \infty$ [3,4]. If the parameter $x$ can be connected to an observable, Eq. (1) corresponds to the Heisenberg uncertainty relation for conjugate variables [1,2]. (Note that we can also exploit quantum “tricks” such as entanglement and squeezing in optimizing the state preparation of the probe and/or the detection stage [17].)

Here we will derive a bound in terms of the expectation value $(\langle H \rangle)$ of $H$, which in the simple case of constant $\Delta X$, takes the form (see Fig. 1)

$$\Delta X \geq \kappa/[(\nu(\langle H \rangle - E_0))] \quad (2)$$

where $E_0$ is the value of a “ground state,” the minimum eigenvalue of $H$ whose eigenvector is populated in the

FIG. 1 (color online). Lower bounds to the precision estimation $\Delta X$ as a function of the experimental repetitions $\nu$. The gray area in the graph represents the forbidden values due to our bound (2). The dashed (dashed-line) area represents the forbidden values due to the Cramér-Rao bound, or the Heisenberg uncertainty, (1). Possible estimation strategies have precision $\Delta X$ that cannot penetrate in the colored regions. For large $\nu$ the Cramér-Rao bound (which scales as $1/\sqrt{\nu}$) is stronger, as expected since in this regime it is achievable. Our bound is not achievable in general, so that the gray area may be expanded when considering specific estimation strategies. [Here we used $(\langle H \rangle - E_0) = 0.1$ (a.u.) and $\Delta H = 4$ (a.u.).]
the joint states $\rho_x$ and $\rho_{x'}$, and where $\alpha(F) \approx \beta^2(F) = 4\arctan^2(\sqrt{F})/\pi^2$ are decreasing functions defined in [21]. The inequality (4) tells us that the parameter difference $|x' - x|$ induced by a transformation $e^{-i(\xi - \eta)H}$ that employs resources $\langle H \rangle - E_0$ and $\Delta H$ cannot be arbitrarily small (when the parameter $x$ coincides with the evolution time, this sets a limit to the "speed" of the evolution, the quantum speed limit).

We now give the main ideas of the proof of (2) by focusing on a simplified scenario, assuming pure probe states $|\psi_{\xi} \rangle = U_{\xi}|\psi \rangle$, and unbiased estimation strategies constructed in terms of projective measurements with RMSE $\Delta X$ that do not depend on $x$. The detailed proof is given in [27], where these restrictions are dropped. For unbiased estimation, $x = \sum_j P_j(x) x_j$ and the RMSE coincides with the variance of the distribution $P_j(x)$, i.e., $\Delta X = \sqrt{\sum_j P_j(x)[x_j - x]^2}$, where $P_j(x) = |\langle x_j | \psi_{\xi} \rangle|^2$ is the probability of obtaining the result $x_j$ while measuring the joint state $|\psi_{\xi} \rangle$ with a projective measurement on the joint basis $|x_j \rangle$. Let us consider two values $x$ and $x'$ of the parameter that are further apart than the measurement's RMSE, i.e., $x' - x = 2\Delta X$ for a $\lambda \approx 1$ that will be specified later. If no such $x$ and $x'$ exist, the estimation is extremely poor: basically the whole domain of the parameter is smaller than (or of the same order of) the RMSE. Hence, we can always assume that such a choice is possible for estimation strategies that are sufficiently accurate to be of interest, as discussed in detail below. The Chebyshev inequality states that for an arbitrary probability distribution $p$, the probability that a result $x$ lies more than $\lambda \Delta X$ away from the average $\mu$ is upper bounded by $1/\lambda^2$, namely $p(|x - \mu| \geq \lambda \Delta X) \leq 1/\lambda^2$. It implies that the probability that measuring $|\Psi_{\xi} \rangle := |\psi_{\xi} \rangle^{\ast \ast}$, the outcome $x_j$ lies within $\lambda \Delta X$ of the mean value associated with $|\Psi_{\xi} \rangle := |\psi_{\xi} \rangle^{\ast \ast}$ cannot be larger than $1/\lambda^2$. By the same reasoning, the probability that measuring $|\Psi_{\eta} \rangle$ the outcome $x_j$ will lie within $\lambda \Delta X$ of the mean value associated with $|\Psi_{\tau} \rangle$ cannot be larger than $1/\lambda^2$. This implies that the overlap between the states $|\Psi_{\xi} \rangle$ and $|\Psi_{\eta} \rangle$ cannot be too large: more precisely $F = |\langle \Psi_{\xi} | \Psi_{\eta} \rangle|^2 \leq 4/\lambda^2$. Replacing this expression into (4) (exploiting the fact that $\alpha$ and $\beta$ are decreasing functions) we obtain

$$2\lambda \Delta X \geq \frac{2}{\rho^2} \max \left[ \frac{\alpha(\lambda^2)}{\rho(\langle H \rangle - E_0)}, \frac{\beta(\lambda^2)}{\sqrt{\rho} \Delta H} \right],$$

whence we obtain (2) by optimizing over $\lambda$ the first term of the max, i.e., choosing $\kappa = \sup_{\lambda} \rho \alpha(4/\lambda^2)/(4\lambda) \approx 0.091$, maximized for $\lambda \approx 4.09$. The second term of the max gives rise to a quantum Cramér-Rao type uncertainty relation (or a Heisenberg uncertainty relation) that, consistently with the optimality of Eq. (1) for $\rho \gg 1$, has a pre-factor $\pi\beta(4/\lambda^2)/(4\lambda)$ that is smaller than $1/2$ for all $\lambda$. This means that for large $\rho$ the bound (2) will be asymptotically
superseded by the Cramér-Rao part, which scales as $\frac{1}{\sqrt{\nu}}$ and is achievable in this regime. In other words, when it is physically significant, the Cramér-Rao bound always wins over the bound (2).

Analogous results can be obtained (see [27]) when considering more general scenarios where the input states of the probes are not pure, the estimation process is biased, and it is performed with arbitrary (possibly adaptive) positive operator-valued measure (POVM) measurements. [In the case of biased measurements, the constant $\kappa$ in (2) and (3) must be replaced by $\kappa = \sup \| a \| (4 / \lambda^2) / [4(\lambda + 1)] \approx 0.074$ (maximized for $\lambda \approx 4.64$) where $a + 1$ term appears in the denominator.] In this generalized context, whenever the RMSE depends explicitly on the value $x$ of the parameter, the result (2) derived above is replaced by the weaker relation (3). Such inequality clearly does not necessarily exclude the possibility that at a “sweet spot” the estimation might violate the scaling (2), as happens, e.g., in the strategies of [14–15,28], which are, hence, fully compatible with our bounds. However, Eq. (3) is still sufficiently strong to exclude accuracies of the form $\Delta X(x) = 1 / R(x, \nu(H))$ where, as in Refs. [15,28], $R(x, z)$ is a function of $z$ that, for all $x$, increases more than linearly, i.e., $\lim_{z \to \infty} R(x, z) = 0$.

The bound (2) has been derived under the explicit assumption that $x$ and $x'$ exist such that $x' - x = 2\Delta X$ for some $\lambda \approx 4.09$, which requires one to have $x' - x \geq 2\Delta X$. This means that the estimation strategy must be good enough: the probe is sufficiently sensitive to the transformation $U_\lambda$ that it is shifted by more than $\Delta X$ during the interaction. The existence of pathological estimation strategies which violate such condition cannot be excluded a priori. Indeed trivial examples of this sort can be easily constructed, a fact which may explain the complicated history of the Heisenberg bound with claims [8–13] and counterclaims [14–16,28]. It should be stressed, however, that the assumption $x' - x \geq 2\Delta X$ is always satisfied except for extremely poor estimation strategies with such large errors as to be practically useless. One may think of repeating such a poor estimation strategy $\nu > 1$ times and of performing a statistical average to decrease its error. However, for sufficiently large $\nu$ the error will decrease to the point in which the $\nu$ repetitions of the poor strategy are, collectively, a good strategy and hence, again subject to our bounds (2) and (3).

Our findings are particularly relevant in the field of quantum optics, where a controversial and long-debated problem [8–16,28] is to determine the scaling of the ultimate limit in the interferometric precision of estimating a phase as a function of the total average energy devoted to preparing the $\nu$ copies of the probes: it has been conjectured [8–13] that the phase RMSE is lower bounded by the inverse of the total number of photons employed in the experiment, the “Heisenberg bound” for interferometry [29]. It corresponds to an equation of the form of Eq. (2), choosing $x = \phi$ (the relative phase between the modes in the interferometer) and $H = a^\dagger a$ (the number operator).

Scalings of this sort have been established for some redefinitions of the uncertainty measure [3] or for specific detection strategies [30] (see, e.g., Refs. [31–33] and references therein) while its achievability for the RMSE measure has been recently proven in [34]. Still its general validity for the RMSE has been questioned several times [14–16,28]. In particular schemes have been proposed [15,28] that apparently permit better scalings in the achievable RMSE (for instance $\Delta X \approx (\nu(H))^{-\gamma}$ with $\gamma > 1$). None of these protocols have conclusively proved such scalings for arbitrary values of the parameter $x$, but a sound, clear argument against the possibility of breaking the $\gamma = 1$ scaling of Eq. (2) was missing up to now. Our results validate the Heisenberg bound by showing that it applies to all those estimation strategies whose RMSE $\Delta X$ does not depend on the value of the parameter $x$ and that the remaining strategies can have good precision only for isolated values of the unknown parameter $x$.

After the appearance of the first version of our manuscript, related papers have appeared in which the bounds derived here are analyzed in the presence of nontrivial prior information [20,35]. Moreover, for optical interferometry our findings have been strengthened in Ref. [36] and a compatible bound dependent on the prior information was obtained through rate distortion theory in [37].

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[1] W. Heisenberg, Z. Phys. 43, 172 (1927); English translation in Quantum Theory and Measurement, edited by J.A. Wheeler and H. Zurek (Princeton Univ., Princeton, NJ, 1983), pp. 62–84.
[2] H. P. Robertson, Phys. Rev. 34, 163 (1929).
[3] A. S. Holevo, Probabilistic and Statistical Aspect of Quantum Theory (Edizioni della Normale, Pisa, 2011).
[4] C. W. Helstrom, Quantum Detection and Estimation Theory (Academic Press, New York, 1976).
[5] S.L. Braunstein, M. C. Caves, and G. J. Milburn, Ann. Phys. (N.Y.) 247, 135 (1996).
[6] S.L. Braunstein and C.M. Caves, Phys. Rev. Lett. 72, 3439 (1994).
[7] V. Giovannetti, S. Lloyd, and L. Maccone, Nature Photon. 5, 222 (2011).
[8] S.L. Braunstein, A. S. Lane, and C.M. Caves, Phys. Rev. Lett. 69, 2153 (1992).
[9] B. Yurke, S. L. McCall, and J. R. Klauder, Phys. Rev. A 33, 4033 (1986).
[10] B. C. Sanders and G. J. Milburn, Phys. Rev. Lett. 75, 2944 (1995).
[11] Z. Y. Ou, Phys. Rev. A 55, 2598 (1997); Z. Y. Ou, Phys. Rev. Lett. 77, 2352 (1996).
[12] J. J. Bollinger, W. M. Itano, D. J. Wineland, and D. J. Heinzen, Phys. Rev. A 54, R4649 (1996).
[13] P. Hyllus, L. Pezzè, and A. Smerzi, Phys. Rev. Lett. 105, 120501 (2010).
[14] P. M. Anisimov G. M. Raterman, A. Chiruvelli, W. N. Plick, S. D. Huver, H. Lee, and J. P. Dowling, Phys. Rev. Lett. 104, 103602 (2010).
[15] A. Rivas and A. Luis, arXiv:1105.6310v2.
[16] Y. R. Zhang et al., arXiv:1105.2990v2.
[17] V. Giovannetti, S. Lloyd, and L. Maccone, Phys. Rev. Lett. 96, 010401 (2006).
[18] N. Margolus, arXiv:1109.4994v2.
[19] Poor strategies that do not increase knowledge beyond the prior information can beat the bound (2), but they are of little practical relevance, see also [20].
[20] V. Giovannetti and L. Maccone, Phys. Rev. Lett. 108, 210404 (2012).
[21] N. Margolus and L. B. Levitin, Physica (Amsterdam) 120D, 188 (1998).
[22] K. Bhattacharyya, J. Phys. A 16, 2993 (1983).
[23] L. Mandelstam and I. G. Tamm, J. Phys. (USSR) 9, 249 (1945).
[24] The fidelity between two states $\rho$ and $\sigma$ is defined as $F = \text{Tr}(\sqrt{\sqrt{\rho} \sigma \sqrt{\rho}})^2$.
[25] A connection between quantum metrology and the Margolus-Levitin theorem was proposed in M. Zwierz, C. A. Pérez-Delgado, and P. Kok, Phys. Rev. Lett. 105, 180402 (2010); but this claim was subsequently retracted by the same authors in M. Zwierz, C. A. Pérez-Delgado, and P. Kok, Phys. Rev. Lett. 107, 059904(E) (2011).
[26] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.108.260405 for the extended proof that covers the case in which the probes are not pure, the estimation process is biased, and it is performed with arbitrary global and local (possibly adaptive) POVM measurements.
[27] J. H. Shapiro, S. R. Shepard, and N. C. Wong, Phys. Rev. Lett. 62, 2377 (1989).
[28] This “Heisenberg” bound [8–13] should not be confused with the Heisenberg scaling defined for general quantum estimation problem [7] in which the $\sqrt{\rho}$ at the denominator of Eq. (1) is replaced by $\rho$ by feeding the $\rho$ inputs with entangled input states—e.g. see Refs. [7,17].
[29] G. S. Summy and D. T. Pegg, Opt. Commun. 77, 75 (1990).
[30] D. W. Berry, B. L. Higgins, S. D. Bartlett, M. W. Mitchell, G. J. Pryde, and H. M. Wiseman, Phys. Rev. A 80, 052114 (2009).
[31] D. W. Berry and H. M. Wiseman, Phys. Rev. Lett. 85, 5098 (2000).
[32] D. W. Berry, H. M. Wiseman, and J. K. Breslin, Phys. Rev. A 63, 053804 (2001).
[33] M. Hayashi, Progr. Informat. 8, 81 (2011), arXiv:1011.2546v1.
[34] M. Tsang, Phys. Rev. Lett. 108, 230401 (2012).
[35] M. J. W. Hall, D. W. Berry, M. Zwierz, and H. W. Wiseman, Phys. Rev. A 85, 041802 (2012).
[36] R. Nair, arXiv:1204.3761v1.