We propose a criterion to classify and compare sources of single photons: the $p$-norm of their $k$th-order photon correlations. While ideally one should use the uniform norm with $p \to \infty$, in practice, the $3$-norm already allows a meaningful comparison and shows the serious limitations of the usual criterion based on the second order correlation $g^{(2)}$ only. We apply this new criterion to a large family of popular single photon sources and show that some of them do not, in fact, qualify for this function. The best sources from our selections consist of a cascade of two-level systems, improving on the otherwise standard low-power resonant excitation of a single two-level system.

Single photon sources (SPS) [1] constitute a basic component of the future quantum technologies [2] and their laboratory implementation is therefore actively pursued by several groups that compete for the delivery of the best emitters with the required characteristics [3–7]. The desiderata for a good SPS cover a large range of both fundamental (antibunching, indistinguishability, etc.) and applied (efficiency, price, etc.) interest [3–10]. Since the technology is yet largely within a development phase, the most prized characteristics are still those of a fundamental character, and there is a need to compare different emitters from various platforms and operating in different conditions. One of the important requisites is the suppression of multiple-photon events, which is usually characterized by the zero-delay second-order correlator function $g^{(2)} = \langle \varsigma \varsigma^\dagger \varsigma \varsigma^\dagger \rangle / \langle \varsigma \varsigma^\dagger \rangle^2$ that quantifies the correlations between two quanta from a field with annihilation operator $\varsigma$ [11]. It has long been known that $g^{(2)} < 1$, linked to sub-Poissonian fluctuations, is a proof of the non-classical character of the field and is commonly used as a criterion for the quantum character of a state. Such a condition accommodates as well, however, a weak-class of quantum states that consists of convex mixtures of Gaussian states [12], while SPS are pursued for providing stronger quantum resources that allow, for instance, the reconstruction of any other state. A single photon, being a very particular case, requires a dedicated way to qualify it. The criterion is simple enough but is not practical, it reads: $g^{(2)} = 0$. However, a mathematical zero is not a good criterion for a physical observable, that is prone to error and approximations. Aiming instead for this quantity to be as close to zero as possible would be the natural remedy but is in fact not adequate, as we discuss in this Letter. Consequently, we propose a fitting substitute and use it to establish a first comparison between different mechanisms that power SPS.

The fundamental emitter of single photons is the two-level system (2LS) with annihilation operator $\varsigma$. By construction, the 2LS can only support one excitation, so it can only emit one photon at a time. The $g^{(2)}$ of its emitted light is, theoretically, exactly zero. This is one example of a mathematical zero since the cancellation of $g^{(2)}$ when $\varsigma = \varsigma^\dagger$ occurs at the level of the nilpotent operator: $\sigma^2 = 0$. This is therefore realized for any excitation of the system. On the detection side, however, the situation is not equally ideal. In fact, already the strong quantum character of the state is affected by the contribution of the vacuum. This has been recently shown [12] through an upper bound for the probability $p(1)$ of the one-photon component for the class of convex mixtures of Gaussian states: no such state can ever be found with this probability overcoming $3\sqrt{3}/(4e) \approx 0.477889$ [12]. This provides a criterion for the strong quantum character of the state. For a 2LS with decay rate $\gamma_\varsigma$ that is excited incoherently at the rate $P_\varsigma$, this is achieved when $P_\varsigma/\gamma_\varsigma > p(1)/(1 - p(1)) \approx 0.915303$. There is therefore a threshold in pumping to meet this criterion. Higher pumping of course satisfy it even more and in the limit of infinite pumping, the incoherent excitation of a 2LS maintains the pure Fock state $|1\rangle$ in the steady state, since the system is forced to remain excited. This is the most quantum state one can get with respect to a single-photon objective. Strong pumping is not, however, a good approach to SPS, since it spoils another important quality: their indistinguishability. With pumping comes power broadening and photons are emitted with increasingly fluctuating energies. The Fock state $|1\rangle$ emits photons one at a time but at all the frequencies!

The photon frequencies are in fact a crucial component when it comes to suppressing multiple-photon emission. In practice, there are numerous reasons to get accidental coincidences, but one is of a fundamental character, and is related to the tail events in the power spectrum (the energy distribution of the emitted photons). This is fundamental because the power spectrum for spontaneous emission of a 2LS is Lorentzian which means that it fluctuates to all orders and it is therefore fundamentally impossible to detect all the photons (since it is impossible to detect at all the frequencies). In contrast, a distribution with a standard deviation would allow, for all practical purposes, to detect all the photons given a wide enough filter (say five standard deviations). Interposing a filter
between the emitter and the detector (this also describes the detector’s finite bandwidth) affects the statistics of the detected photons \[13\]. Thus, although the 2LS can host only one excitation, its filtering results in nonzero probabilities \( p(n) \) to detect \( n \geq 2 \) photons. There are various ways to understand this result: the filter introduces a time uncertainty that makes it possible to pile-up consecutive photons. From a quantum mechanical viewpoint, the 2LS emits a continuous stream of photons at all the possible frequencies, which amplitudes interfere destructively to result in a single photon upon detection. Trimming the tails spoils the destructive interference and allows more than one photon to reach the detector.

The distributions for an incoherently excited 2LS that is detected in various frequency windows is shown as joined dots in Fig. 1(a). Here the pumping was very large to keep the system in its excited state. This is obtained through a master equation for the incoherent excitation (at a rate \( P_\sigma \)) of a 2LS with free Hamiltonian \( H_\sigma = \omega_\sigma \sigma^\dagger \sigma \) and decay rate \( \gamma_\sigma \) (we take \( \hbar = 1 \) along the text):

\[ \partial_t \rho = i[\rho, H_\sigma] + \frac{\gamma_\sigma}{2} \mathcal{L}_\sigma \rho + \frac{P_\sigma}{2} \mathcal{L}_\sigma^\dagger \rho, \]

where \( \mathcal{L}_c \rho = (2\epsilon \rho c^\dagger - c^\dagger c \rho - \rho c^\dagger c) \). The filtering, or the finite temporal and spectral resolution of the detectors, are taken into account self-consistently through the theory of frequency-filtered and time-resolved \( N \)-photon correlations \[13\], where a sensor with spectral width \( \Gamma \) (and hence with temporal resolution \( 1/\Gamma \)) is included into the dynamics of the system in the limit of its vanishing coupling to the SPS, so that it only “senses” correlations without perturbing the system. For an incoherently driven 2LS, the \( n \)-th order filtered correlation function \( g_\Gamma^{(n)} \) can be thus obtained recursively using the relation:

\[ g_\Gamma^{(n)} = g_\Gamma^{(n-1)} \frac{n \Gamma_\sigma (\Gamma_\sigma + \Gamma)}{(\Gamma_\sigma + (n-1) \Gamma) [\Gamma_\sigma + (2n-1) \Gamma]}, \]

where \( g_\Gamma^{(1)} = 1 \) and \( \Gamma_\sigma \equiv \gamma_\sigma + P_\sigma \). Since the filtered \( n \)-photon correlations are nonzero, the effective density matrix that describes the state of the 2LS is clearly no longer constrained to one excitation. The probability \( p(n) \) to find in the Fock state \( |n \rangle \) the effective emitter that corresponds to a filtered 2LS can be recovered through the bijective mapping provided in Ref. \[14\]:

\[ p(n) = \sum_{k \geq n} \frac{(-1)^{k+n}}{n! (k-n)!} G_\Gamma^{(k)}, \]

where \( G_\Gamma^{(k)} = \langle \sigma^\dagger \sigma \rangle^k G_\Gamma^{(k)} \) is the \( k \)-th order unnormalized correlator, and \( \langle \sigma^\dagger \sigma \rangle_\gamma = P_\sigma \Gamma / (\Gamma_\sigma \Gamma + \Gamma^2_\sigma) \) is the population that passes through a Lorentzian bandpass filter. Combining Eqs. \[20\], we obtain the distribution of the filtered SPS in closed form:

\[ p(n) = C_n \times _2F_2 \left( n + 1; \frac{2n + 1}{2} + \frac{\Gamma_\sigma}{\Gamma}, n + \frac{\Gamma_\sigma}{\Gamma} + \frac{P_\sigma}{2\Gamma}; -\frac{P_\sigma}{2\Gamma} \right), \]

where \( _2F_2 \) is the generalized hypergeometric function and \( C_n \) are coefficients that satisfy:

\[ C_n - C_{n-1} = \frac{P_\sigma \Gamma}{[\Gamma_\sigma + (n-1) \Gamma] [\Gamma_\sigma + (2n-1) \Gamma]} \]

with \( C_0 = 1 \).

These expressions, displayed in Fig. 1(a), show that filtering (or, equivalently, using a low-resolution detector) spoils both \( g(2) \) and the strong quantum character of the source, as it brings it to break the Filip & Mista aforementioned criterion \[12\] already for a filter (or detection) bandwidth \( \Gamma \) commensurate to the power broadened linewidth \( \Gamma_\sigma \) (we also show as a bar plot the
probability distribution for the Gaussian state that maximizes the probability of one excitation $p(1)$. Narrower filters completely thermalize the signal, with $g^{(n)} \to n!$ when $\Gamma/\Gamma_\sigma \to 0$. Figures 1(b–c) show the probability for the vacuum $p(0)$ and for one-excitation $p(1)$ as a function of the incoherent pumping rate and filtering of the 2LS. Taken together, these two quantities provide a stricter criterion to separate weak classes of quantum states from strong ones [15]. Namely, the parameter-space that lies on the right of the thick blue line realizes steady state of the 2LS that cannot be represented as a convex mixture of Gaussian states, i.e., are the quantum non-Gaussian states sought for quantum applications. The processes shaping such states require high-order quantum nonlinearities. This confirms that to maintain the quantum character of the source, one requires a minimum filter linewidth. For a pumping rate commensurable with the decay rate, such a minimum is roughly twice the intrinsic linewidth of the 2LS. As could be expected, for higher pumping rates the power-broadening implies that a correspondingly larger filter is required. However, at vanishing pumping, this time unexpectedly, one also needs increasing filtering, showing how the more scarcely the 2LS is excited, the less quantum it is. This comes from the deleterious contribution of the vacuum.

Since the tail events are the main responsible for spoiling the SPS, an emitter whose power spectrum is not Lorentzian but has faster decaying tails would allow to collect more of the photons in a reduced spectral window and thus better preserve its quantum features. The simplest way to achieve that is to turn to coherent driving, in which case Heitler processes of coherent absorption and emissions trim the Lorentzian tails to yield instead a Student $t_2$ distribution, which still has fat tails but that decrease faster than a Lorentzian [16]. The same idea can be further enforced by turning to quantum rather than classical excitation, in a cascaded scheme of several 2LS in a chain [15]. A drawback of coherent driving is that it forbids total inversion of the 2LS, which, due to stimulated emission, is at most half-excited. Therefore, while $g^{(2)}$ is indeed smaller, the quantum state itself deviates less from classical mixtures than under incoherent excitation. This is shown in Figs. 1(d–e) that are the counterpart of the above row but for coherent excitation. The probability saturates at $p(1) = 1/2$, which is a considerable impediment when trying to demonstrate the genuine quantum character of a source (for instance through the measurement of a negative Wigner function).

From our discussion so far, it is clear that the quantum non-Gaussian character of the state, although of great fundamental interest, is not the best practical way to assess the quality of a SPS. The widely used criterion that rely on $g^{(2)}$ is more adequate and indeed sets the standard. It presents, however, some subtleties and is frequently mishandled, for instance through a popular criterion $g^{(2)} < 1/2$ that supposedly guarantees the single-photon or single-emitter character of the source [17–21]. We have shown in an earlier work [15, 22] however, that the Hilbert space contains a myriad of states with population larger than one that satisfy the condition $g^{(2)} < 1/2$. The exact criterion remains $g^{(2)} = 0$. It is also invalid, however, to aim for $g^{(2)} \to 0$, as states with arbitrarily small but nonzero $g^{(2)}$, but featuring a huge $g^{(n)}$ for $n \geq 3$, are possible. The state $\rho = (297001/300000)|0\rangle\langle 0| + (1999/200000) |1\rangle\langle 1| + (1/600000) |3\rangle\langle 3|$ for instance has $g^{(2)} = 1/10$ and $g^{(3)} = 10$, so stopping at the second-order correlation does not well characterize a SPS. Clearly, in the most general setting, a SPS has to exclude all fanciful designs, possibly deviously conceived fake SPS, that would give the illusion of good antibunching through standard $g^{(2)}$ measurements, while opening photon-number splitting backdoors for hacking cryptographic protocols [23] by encoding the information in $N$ photon states not detectable at the two-photon level.

Here, we propose a novel criterion to well identify and compare usefully SPS, by using all the correlation functions. Namely, we propose the $N$-norm of the vector of correlation functions:

$$\|(g^{(k)})\|_N \equiv \sqrt{\sum_{k=2}^{N+1} (g^{(k)})^N},$$

that quantifies the deviation between the ideal SPS, which is perfectly antibunched to all orders, and an actual source. The criterion is then that $\|(g^{(k)})\|_N$ should be as small as possible (not simply the $g^{(2)}$, which is the 1-norm) for $N$ as high as technically feasible. In the limit $N \to \infty$, the quantity becomes $\|(g^{(k)})\|_\infty = \max(g^{(k)})$ and measures the failure of the SPS by the order at which it less suppresses multiple-photon emission.

In Fig. 2(a), we show the result for several popular sources, limiting to the 3-norm (i.e., measuring up to $g^{(4)}$), which is enough for the cases we have chosen (that do not engineer scenarios where the SPS character breaks for a large number of photons). We show the results as a function of the filtering spectral resolution, that, as we have already discussed, is a key parameter for photon correlations. The incoherently driven 2LS sets a useful reference. As the filtering window is enlarged, our measure [3] indeed decreases. The coherently driven 2LS, in comparison, decreases much faster, and this confirms the known fact that resonant excitation provides better antibunching than the non-resonant case. This is the case at low driving, since at high driving, the system enters the Mollow triplet regime [24], in which case the statistics of the emission varies strongly with the frequency [25]. In particular, the emission at frequencies between the triplet’s central and lateral peaks is largely bunched, as opposed to the resonant emission which remains antibunched. Even though performing as a SPS when filtering over the entire structure, this affects con-
by the annihilation operators \( \sigma \) momentum, namely +1 and -1 excitons with opposite third component of the angular pair of excitons (an electron-hole pair). It supports two of a biexciton, i.e., a molecule that consists of a bound that fail in other respects. This the case for instance The linearly polarized states of the biexciton system,\[ g \]

\[ \frac{\gamma}{\gamma} = 40 \text{ and } \Omega /\gamma = 10, \] on the other hand, fails to overcome the benchmark in either polarization. The blockade mechanisms provide better antibunching in the limit of \( \chi/\gamma \sigma = 40 \) and \( \Omega_\sigma/\gamma_\sigma = 10 \), on the other hand, fails to overcome the benchmark in either polarization. The blockade schemes that we have recently proposed [16] and already collected. The situation is even more serious for the unconventional polariton blockade that, contrary to a popular [27] and unconventional [28–30] form. In the former case, antibunching follows from self-interactions (polaritons in this context are essentially strongly interacting photons), while in the latter case, a much stronger antibunching is obtained at vanishing pumping from an interference. Calling \( a \) and \( b \) the annihilation operators of two polariton modes, the Hamiltonian reads\[ H_B = \omega (a^\dagger a + b^\dagger b) + g (a^\dagger b + b^\dagger a) + U (a^\dagger a^\dagger a + b^\dagger b^\dagger b) + \Omega (a^\dagger e^{i\omega t} + ae^{-i\omega t}) \] where \( \omega \) is the polariton energy, \( g \) is the coherent coupling between polaritons (e.g., by tunneling), \( U \) is the strength of polariton interaction, \( \gamma_\sigma \) is the polariton decay rate and \( \Omega a/\sqrt{\gamma_\sigma} \) is the amplitude of the laser driving one of the polaritons. This Hamiltonian covers both types of blockades depending on the frequency of resonant excitation and the coherent coupling. Figure 2(a) shows that the conventional polariton blockade (obtained by setting \( g = 0 \) in \( H_B \) to keep only polaritons \( a \)) behaves differently from the other SPS, as it remains roughly constant as more light is collected. The situation is even more serious for the unconventional polariton blockade that, contrary to a popular belief that made its manifestation actively sought after, performs rather poorly as a SPS, in particular in the configuration where its \( g(2) \) is very small. Indeed, despite its excellent antibunched \( g(2) \) emission, it also exhibits a superbunched \( g(k) \) for \( k \geq 3 \), and therefore fails to behave as a SPS for most of the values of \( U/\gamma_\sigma \), including that where \( g(2) \ll 1 \). Such a discrimination between the two types of blockades is shown in Fig. 2(b), comparing

![FIG. 2. (Color online). Criterion for Single Photon Sources.](image)

(a) 3-norm for several sources of light as a function of their spectral filtering. The incoherently driven 2LS (in the limit of \( P_\sigma \ll \gamma_\sigma \)) is set as the benchmark for SPS. It is always surpassed by the coherently driven 2LS and the cascaded 2LS as long as \( \Omega_\sigma \ll \gamma_\sigma \). The emission of the biexciton (with \( \chi/\gamma_\sigma = 40 \) and \( \Omega_\sigma/\gamma_\sigma = 10 \)), on the other hand, fails to overcome the benchmark in either polarization. The blockade schemes that we have recently proposed [16] and already discussed above is confirmed as a better SPS than under classical resonant driving, although the slope is the same as for the conventional case. It is in fact the best SPS of our selection.

A value of our criterion resides in its identification of sources with a strong antibunching (small \( g(2) \)) but that fail in other respects. This the case for instance of a biexciton, i.e., a molecule that consists of a bound pair of excitons (an electron-hole pair). It supports two excitons with opposite third component of the angular momentum, namely +1 and -1, which are represented by the annihilation operators \( \sigma_\uparrow \) and \( \sigma_\downarrow \), respectively. The linearly polarized states of the biexciton system, namely \( \sigma_{H,V} = (\sigma_\uparrow + \sigma_\downarrow)/\sqrt{2} \), describe transitions from the exciton to the ground state emitting a vertically and horizontally polarized photon, respectively. To test the biexciton as a SPS, we drive it coherently with a vertically polarized laser of amplitude \( \Omega_\sigma/\sqrt{\gamma_\sigma} \) and in resonance with the exciton transition. The Hamiltonian describing the biexciton is given by\[ H_B = \omega_\sigma (\sigma_\uparrow^\dagger \sigma_\uparrow + \sigma_\downarrow^\dagger \sigma_\downarrow) - \chi (\sigma_\uparrow^\dagger \sigma_\downarrow^\dagger \sigma_\downarrow \sigma_\uparrow + \Omega_\sigma (\sigma_\uparrow e^{i\omega t} + \sigma_\downarrow e^{-i\omega t}) \] where \( \omega_\sigma \) is the bare exciton energy (we are considering resonant excitons) and \( \chi \) is the bonding energy, so that the biexciton energy is \( 2\omega_\sigma - \chi \). The biexciton is described by the master equation given in Eq. (1) replacing \( H_\sigma \) by \( H_B \), dropping the term related to the incoherent driving \((L_\sigma, \rho)\), and letting the excitons decay with rate \( \gamma_\sigma \) for both polarizations. Figure 2(a) shows that the 3-norm for the emission of the biexciton saturates above the measure set by the incoherently driven 2LS. This means that, as expected, the biexciton is not a good SPS. However, Fig. 2(a) also shows that the asymmetry in the driving field favors the vertically polarized emission, which performs better than the horizontally polarized emission. In fact, for a large region of values for \( \Gamma/\gamma_\sigma \), the vertically polarized emission of the biexciton is a better SPS than the central peak of the Mollow triplet.

Another striking case is the so-called “polariton blockade” (also known as “photon blockade”), in both its conventional [27] and unconventional [28–30] form. In the former case, antibunching follows from self-interactions (polaritons in this context are essentially strongly interacting photons), while in the latter case, a much stronger antibunching is obtained at vanishing pumping from an interference. Calling \( a \) and \( b \) the annihilation operators of two polariton modes, the Hamiltonian reads\[ H_B = \omega (a^\dagger a + b^\dagger b) + g (a^\dagger b + b^\dagger a) + U (a^\dagger a^\dagger a + b^\dagger b^\dagger b) + \Omega (a^\dagger e^{i\omega t} + ae^{-i\omega t}) \] where \( \omega \) is the polariton energy, \( g \) is the coherent coupling between polaritons (e.g., by tunneling), \( U \) is the strength of polariton interaction, \( \gamma_\sigma \) is the polariton decay rate and \( \Omega a/\sqrt{\gamma_\sigma} \) is the amplitude of the laser driving one of the polaritons. This Hamiltonian covers both types of blockades depending on the frequency of resonant excitation and the coherent coupling. Figure 2(a) shows that the conventional polariton blockade (obtained by setting \( g = 0 \) in \( H_B \) to keep only polaritons \( a \)) behaves differently from the other SPS, as it remains roughly constant as more light is collected. The situation is even more serious for the unconventional polariton blockade that, contrary to a popular belief that made its manifestation actively sought after, performs rather poorly as a SPS, in particular in the configuration where its \( g(2) \) is very small. Indeed, despite its excellent antibunched \( g(2) \) emission, it also exhibits a superbunched \( g(k) \) for \( k \geq 3 \), and therefore fails to behave as a SPS for most of the values of \( U/\gamma_\sigma \), including that where \( g(2) \ll 1 \). Such a discrimination between the two types of blockades is shown in Fig. 2(b), comparing
all $\|g^{(k)}\|_p$ for $p = 1, \ldots, 4$ and this time as a function of the interaction strength (normalized to the broadening) $U/\gamma_a$. The dashed lines correspond to the unconventional blockade and show an essentially converged result regardless of $p$, indicating that the conventional blockade is a SPS that gets increasingly (albeit slowly) better with increasing nonlinearity. On the other hand, while the unconventional blockade is much better as perceived through the 1-norm (the conventional $g^{(2)}$, this is the result hailed in the literature), it is quickly spoiled at the 2-norm level, particularly at the value of $U/\gamma_a$ that minimizes $g^{(2)}$. It remains a better SPS than its conventional blockade counterpart on the range shown here (dashed line), but in a different configuration than the one that is aimed for (dotted line) and would be overtaken by the conventional blockade for stronger nonlinearities.

With such a criterion in hand, one can devise configurations that combine the assets of various approaches. For instance, the bare 2LS presents the shortcoming of emitting in a large solid angle and placing it in a cavity brings several advantages, such as increasing its emission rate by Purcell enhancement, but also for practical purposes, in collecting and directing the light in a focused output beam. One can ask whether this has a cost in other aspects and weakens the supression of multiple-photon emission as compared to bare SPS. We find with Eq. (5) that in the limit of weak-coupling and fast decay rate of the cavity, the impact is negligible, which is a welcome result as this configuration is widespread. More involved designs can optimize the combination of these aspects in other configurations, for instance the detuned photon blockade [31] proves to remain an excellent SPS within a cavity, making such an implementation a serious contender for future applications.

In summary, we studied the quantum non-Gaussian character of a 2LS and how filtering and/or finite-detector resolution is affecting its quality as a SPS. We proposed a new criterion to quantify its performance: the $N$-norm of the photon correlation functions. For practical purposes, it is convenient and often enough to limit to the 2-norm or 3-norm. On the other hand, we showed that $g^{(2)}$ alone (the 1-norm), which is the widespread standard, is not sufficient and can lead to erroneous assessments. We illustrated the criterion with a variety of mechanisms to implement SPS and found that the yet-to-be-implemented cascaded SPS provide the best realization of single-photon emitters while resonance fluorescence at low driving are the best already existing SPS. We also showed how other sources that appear to be good SPS do actually fail in a stricter sense that could jeopardize quantum information protocols.

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[1] Aharonovich, I., Englund, D. & Toth, M. Solid-state single-photon emitters. Nat. Photon. 10, 631 (2016). URL doi:10.1038/nphoton.2016.186
[2] O’Brien, J. L., Furusawa, A. & Vuckovic, J. Photonic quantum technologies. Nat. Phys. 3, 687 (2009). URL doi:10.1038/nphys687
[3] Kühnmann, A. V. et al. Transform-limited single photons from a single quantum dot. Nat. Commun. 6, 8204 (2015). URL doi:10.1038/ncomms9204
[4] Somaschi, N. et al. Near-optimal single-photon sources in the solid state. Nat. Photon. 10, 340 (2016). URL doi:10.1038/nphoton.2016.23
[5] Ding, X. et al. On-demand single photons with high extraction efficiency and near-unity indistinguishability from a resonantly driven quantum dot in a micropillar. Phys. Rev. Lett. 116, 020401 (2016). URL doi:10.1103/PhysRevLett.116.020401
[6] Wang, H. et al. Near-transform-limited single photons from an efficient solid-state quantum emitter. Phys. Rev. Lett. 116, 213601 (2016). URL doi:10.1103/PhysRevLett.116.213601
[7] Kim, J.-H., Cai, T., Richardson, C. J. K., Leavitt, R. P. & Waks, E. Two-photon interference from a bright single-photon source at telecom wavelengths. Optica 3, 577 (2016). URL doi:10.1364/OPTICA.3.000577
[8] Dada, A. C. et al. Indistinguishable single photons with flexible electronic triggering. Optica 3, 493 (2016). URL 10.1364/OPTICA.3.000493
[9] Müller, K. et al. Self-homodyne-enabled generation of indistinguishable photons. Optica 3, 931 (2016). URL 10.1364/OPTICA.3.000931
[10] Eisaman, M. D., Fan, J., Migdall, A. & Poloyakov, S. V. Single-photon sources and detectors. Rev. Sci. Instrum. 82, 071101 (2011). URL doi:10.1063/1.3610677
[11] Paul, H. Photon antibunching. Rev. Mod. Phys. 54, 1061 (1982). URL doi:10.1103/RevModPhys.54.1061
[12] Filip, R. & L. Mišta, J. Detecting quantum states with near-optimal single-photon sources from a resonantly driven quantum dot in a micropillar. Phys. Rev. Lett. 107, 213602 (2011). URL doi:10.1103/PhysRevLett.107.213602
[13] Filip, R. & L. Mišta, J. Detecting quantum states with near-optimal single-photon sources from a resonantly driven quantum dot in a micropillar. Phys. Rev. Lett. 107, 213602 (2011). URL doi:10.1103/PhysRevLett.107.213602
[14] Zubizarreta Casalengua, E., López Carreño, J. C., del Valle, E. & Laussy, F. P. Structure of the harmonic oscillator hilbert space. arXiv:1607.03976 (2016).
[15] Ježek, M. et al. Experimental test of the quantum non-gaussian character of a heralded single-photon state. Phys. Rev. Lett. 107, 213602 (2011). URL doi:10.1103/PhysRevLett.107.213602
[16] López Carreño, J. C., Sánchez Muñoz, C., del Valle, E. & Laussy, F. P. Exciting with quantum light. ii. exciting a two-level system. arXiv:1609.00857 (2016).
[17] Michler, P. et al. A quantum dot single-photon turnstile device. Science 290, 2282 (2000). URL doi:10.1126/science.290.5500.2282
[18] Verma, V. B. et al. Photon antibunching from a single lithographically defined InGaAs/GaAs quantum dot. Opt. Express 19, 4182 (2011). URL doi:10.1364/OE.
[19] Martino, G. D. et al. Quantum statistics of surface plasmon polaritons in metallic stripe waveguides. *Nano Lett.* **12**, 2504 (2012). URL doi:10.1021/nl300671w

[20] Lukishova, S. G., Bissell, L. J., Winkler, J., & Stroud, C. R. Resonance in quantum dot fluorescence in a photonic bandgap liquid crystal host. *Opt. Lett.* **37**, 1259 (2012). URL doi:10.1364/OL.37.001259

[21] Palacios-Berraquero, C. et al. Atomically thin quantum light-emitting diodes. *Nat. Comm.* **7**, 12978 (2016). URL doi:10.1038/ncomms12978

[22] López Carreño, J. C. & Laussy, F. P. Exciting with quantum light. i. exciting an harmonic oscillator. *arXiv:1601.06187* (2016).

[23] Brassard, G., Lütkenhaus, N., Mor, T. & Sanders, B. C. Limitations on practical quantum cryptography. *Phys. Rev. Lett.* **85**, 1330 (2000). URL doi:10.1103/PhysRevLett.85.1330

[24] Lemonde, M.-A., Didier, N. & Clerk, A. A. Antibunching and unconventional photon blockade with gaussian squeezed states. *Phys. Rev. A* **90**, 063824 (2014). URL doi:10.1103/PhysRevA.90.063824

[25] Mollow, B. R. Power spectrum of light scattered by two-level systems. *Phys. Rev.* **188**, 1969 (1969). URL doi:10.1103/PhysRev.188.1969

[26] González-Tudela, A., Laussy, F. P., Tejedor, C., Hartmann, M. J. & del Valle, E. Two-photon spectra of quantum emitters. *New J. Phys.* **15**, 033036 (2013). URL doi:10.1088/1367-2630/15/3/033036

[27] Verger, A., Ciuti, C. & Carusotto, I. Polariton quantum blockade in a photonic dot. *Phys. Rev. B* **73**, 193306 (2006). URL doi:10.1103/PhysRevB.73.193306

[28] Liew, T. C. H. & Savona, V. Single photons from coupled quantum modes. *Phys. Rev. Lett.* **104**, 183601 (2010). URL doi:10.1103/PhysRevLett.104.183601

[29] Bamba, M., İmamoğlu, A., Carusotto, I. & Ciuti, C. Origin of strong photon antibunching in weakly nonlinear photonic molecules. *Phys. Rev. A* **83**, 021802(R) (2011). URL doi:10.1103/PhysRevA.83.021802

[30] Flayac, H. & Savona, V. Input-output theory of the unconventional photon blockade. *Phys. Rev. A* **88**, 033836 (2013). URL doi:10.1103/PhysRevA.88.033836

[31] Müller, K. et al. Coherent generation of nonclassical light on chip via detuned photon blockade. *Phys. Rev. Lett.* **114**, 233601 (2015). URL doi:10.1103/PhysRevLett.114.233601