New Journal of Physics
The open access journal at the forefront of physics

Deutsche Physikalische Gesellschaft DPG
IOP Institute of Physics

Published in partnership with: Deutsche Physikalische Gesellschaft and the Institute of Physics

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RECEIVED
6 February 2018

ACCEPTED FOR PUBLICATION
1 March 2018

PUBLISHED
28 March 2018

CORRIGENDUM

Corrigendum: Dynamical modeling of pulsed two-photon interference (2016 *New J. Phys.* 18 113053)

Kevin A Fischer[1], Kai Müller[2], Konstantinos G Lagoudakis[1] and Jelena Vučković[1]

[1] E. L. Ginzton Laboratory, Stanford University, Stanford CA 94305, United States of America
[2] Walter Schottky Institut, Technische Universität München, D-85748 Garching bei München, Germany

E-mail: kevinf@stanford.edu

We were alerted by Chris Gustin that our normalization of the pulse-wise second-order coherence is slightly nonstandard and deserves a deeper explanation. We are very grateful for his help in clarifying equations (44) and (54) in the main text, and we apologize to any readers who may have been confused by this definition. Please note that the clarifications discussed here do not affect the results or conclusions in the text.

First, we briefly describe how to arrive from equations (42) to (44) in the main text. Consider equation (42) in the limit

$$\begin{aligned}&\lim_{k\to\infty} G^{(2)}_{c'd'}(t,\, t'+kt_r)\\ &=\frac{1}{4}\sum_{(e,f)\in\prod(a,b)}[G^{(2)}_{ee}(t,\, t'+kt_r)+G^{(1)}_{e}(t,\, t)G^{(1)}_{f}(t'+kt_r,\, t'+kt_r)\\ &\quad-[G^{(1)}_{e}(t,\, t'+kt_r)]^{*}G^{(1)}_{f}(t,\, t'+kt_r)].\end{aligned} \tag{1}$$

Note that the first-order coherence inherits the envelope of coherence decay from the excitation laser for long times, and hence always vanishes in the long time limit

$$\lim_{k\to\infty} G^{(1)}_{e}(t,\, t'+kt_r)=0. \tag{2}$$

This is an experimental consideration that may be difficult in observing due to various long-time effects such as blinking, potentially limited laser coherence, or operating the correlator in a start-stop configuration (see the main text). In the long time limit, the second-order coherence terms in equation (1) become uncorrelated

$$\lim_{k\to\infty} G^{(2)}_{ee}(t,\, t'+kt_r)=\langle :e^{\dagger}(t)e(t)e^{\dagger}(t'+kt_r)e(t'+kt_r):\rangle \tag{3}$$

$$=\langle e^{\dagger}(t)e(t)\rangle\langle e^{\dagger}(t'+kt_r)e(t'+kt_r)\rangle. \tag{4}$$

Then, using the fact that every pulse period is identical

$$\lim_{k\to\infty} G^{(2)}_{ee}(t,\, t'+kt_r)=G^{(1)}_{e}(t,\, t)G^{(1)}_{e}(t',\, t'), \tag{5}$$

and also the photon flux terms in equation (1) have a similar point

$$\lim_{k\to\infty} G^{(1)}_{f}(t'+kt_r,\, t'+kt_r)=G^{(1)}_{f}(t',\, t'). \tag{6}$$

As a result, we can write

$$\begin{aligned}&\lim_{k\to\infty} G^{(2)}_{c'd'}(t,\, t'+kt_r)\\ &=\frac{1}{4}\sum_{(e,f)\in\prod(a,b)} G^{(1)}_{e}(t,\, t)G^{(1)}_{e}(t',\, t')+G^{(1)}_{e}(t,\, t)G^{(1)}_{f}(t',\, t').\end{aligned} \tag{7}$$

Integrating to pulse-wise form with equation (43) of the paper

$$G^{(2)}_{c'd'}[0]=\int_0^T\int_0^T \mathrm{d}t\mathrm{d}t'\; G^{(2)}_{c'd'}(t,\, t'), \tag{8}$$

and using the definition $\langle\hat{M}_e(T)\rangle=\int_0^T \mathrm{d}t\;\, G^{(1)}_{e}(t,\, t)$, we have the result of equation (44) in the paper

$$\lim_{k\to\infty} G^{(2)}_{c'd'}[kt_r]=\frac{1}{4}(\langle\hat{M}_a(T)\rangle+\langle\hat{M}_b(T)\rangle)^2. \tag{9}$$

The potentially confusing point of this definition is that for a good single-photon source

$$\lim_{k\to\infty} G^{(2)}_{c'd'}[kt_r] \neq G^{(2)}_{c'd'}[t_r] \tag{10}$$

because the first-order coherence terms in equation (1) interfere with the intensity for short times resulting in

$$G^{(2)}_{c'd'}[t_r] \approx \frac{1}{4}(\langle \hat{M}_a(T)\rangle^2 + \langle \hat{M}_b(T)\rangle^2). \tag{11}$$

Hence, using the normalization $G^{(2)}_{c'd'}[0]/G^{(2)}_{c'd'}[t_r]$ results in a denominator that depends on the first-order coherence, which we believe is not the most ideal definition. As Chris pointed out, an alternative way of achieving our preferred normalization is to use the average of the cross- and auto-correlations

$$g^{(2)}_{c'd'}[0] = \lim_{k\to\infty}\frac{G^{(2)}_{c'd'}[0]}{G^{(2)}_{c'd'}[kt_r]} = \frac{4G^{(2)}_{c'd'}[0]}{2G^{(2)}_{c'd'}[t_r] + G^{(2)}_{c'c'}[t_r] + G^{(2)}_{d'd'}[t_r]}, \tag{12}$$

where

$$2G^{(2)}_{c'd'}[t_r] + G^{(2)}_{c'c'}[t_r] + G^{(2)}_{d'd'}[t_r] = \frac{1}{4}(\langle \hat{M}_a(T)\rangle + \langle \hat{M}_b(T)\rangle)^2. \tag{13}$$

Here, we used the fact that $g^{(2)}_{c'd'}[kt_r] = g^{(2)}_{d'c'}[kt_r]$, and assumed that there are no blinking effects. Another common experimental trick to get a normalization by the intensity in equation (11) is to introduce distinguishability (e.g. via polarization rotation) between the two sources so that the fields cannot interfere at the detectors.

We also note a few typos regarding the spontaneous emission rate γ. Occasionally we wrote $1/\gamma$ instead of γ, e.g. in equations (39) and (59) of the main text—the correct expressions are $\hat{g}^{(2)}[0] = \hat{g}^{(2)}_{\text{HOM}}[0] = 0.4\gamma\ \tau_{\text{FWHM}} \pm 0.003$. Similarly, we correct a few of our definitions of τ_{FWHM}, which should read $\tau_{\text{FWHM}} = 0.1/\gamma$ for short pulses or $\tau_{\text{FWHM}} = 3.3/\gamma$ for long pulses.

Finally, we note an error in equation (51) of the main text. We mean to define

$$|g_a^{(1)}[0]|^2 = \frac{|G_a^{(1)}[0]|^2}{\langle \hat{M}_a(T)\rangle^2} \quad \text{and} \quad |G_a^{(1)}[0]|^2 \equiv \int_0^T\int_0^T \mathrm{d}t\ \mathrm{d}t' |G_a^{(1)}(t, t')|^2, \tag{14}$$

which was done correctly in the iPython notebooks detailing our calculational technique. Often, this parameter is referred to as the visibility V. We chose to avoid calling this parameter visibility because its definition does not match the general definition for an arbitrary interferometer. We also comment that for a single-photon source with no error rate, $|g_a^{(1)}[0]|^2$ is precisely the trace purity of the single-photon emission. Given a finite error rate, this equivalence no longer holds. For instance, consider the resonantly driven two-level system with no dephasing: the emitted state is a pure state (with unity trace purity) even though $|g_a^{(1)}[0]|^2 < 1$ (as discussed in the main text).

ORCID iDs

Kevin A Fischer https://orcid.org/0000-0001-9390-1892