

## Realizing the Multiparticle Hanbury Brown–Twiss Interferometer Using Nitrogen-Vacancy Centers in Diamond Crystals

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We demonstrate that the multiparticle Hanbury Brown–Twiss interferometer can be realized in a network of nitrogen-vacancy centers: for an  $N$ -particle system, the interference effect is manifested only in the  $N$ th-order intensity correlation function. The interference effect can be enhanced through a postselection process in which the multipartite Greenberger-Horne-Zeilinger entanglement is generated and tested with Svetlichny inequality.

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The Hanbury Brown–Twiss (HBT) interferometer [1], which was originally devised to determine the angular diameter of stars, shows that phase-sensitive second-order intensity correlations are observed by two spatially separated photodetectors yet not in the intensities of individual photodetectors. HBT effect was subsequently observed in the laboratory by using two independent laser sources [2]. This purely quantum mechanical effect can be attributed to the entanglement that arises from the exchange amplitudes of the two indistinguishable photons and is detected by the violation of Bell inequalities. Recently HBT experiments demonstrating fermionic antibunching have been reported for an electron source [3–6], including theoretical proposals for multiple sources [7,8]. For particles emitted by a thermal source, measurement of second-order correlations, for instance, the fluctuations in intensity-intensity or current-current correlation, yield different results, depending on the effect of particle statistics: positive (negative) for bosonic (fermionic) sources [3]. While for special statistics, bosonic sources can have negative correlation (“antibunching”) [9]. In this Letter, we propose a strategy to realize a multiparticle HBT interferometer in a network of nitrogen-vacancy (NV) color centers in diamond [10].

The network we consider consists of  $N$  chains of NV centers (see Fig. 1 for  $N = 5$ ). Each chain contains two NV centers. The NV center is composed of a nuclear spin  $I = 1/2$  associated with a nitrogen atom  $^{15}\text{N}$  substituting for a carbon atom, and an electronic pair in the spin triplet state  $S = 1$  [11]. The nuclear spins in the  $i$ th chain ( $C_i$ ) are labeled from outside to inside as  $s_{i,1}$  and  $s_{i,2}$ . The electronic spin associated with  $s_{i,j}$  is labeled as  $e_{i,j}$ . The system of NV centers is a promising candidate for quantum information processing, as it can be manipulated at room temperature and the nuclear spin has a long coherence time [12]. The nuclear and electronic spins in a single NV center are coupled through a hyperfine Hamiltonian  $H_{n,e} \sim AI_z S_z$ , where  $A$  is the coupling strength between the nuclear

spin  $I_z$  and the electronic spin  $S_z$ . Two levels ( $m_S = 0, 1$ ) of the electronic spin are utilized as a qubit due to the unequally spaced levels rendered by zero-field splitting [13]. The nuclear spin serves to store quantum information and the electronic pair is used to read out and mediate interactions between nuclear spins of neighboring NV centers. The electronic spin manipulation can be achieved with a microwave-field excitation, and also a controlled- $Z$  gate between the nuclear and electronic spins is realized through the time evolution of the hyperfine Hamiltonian

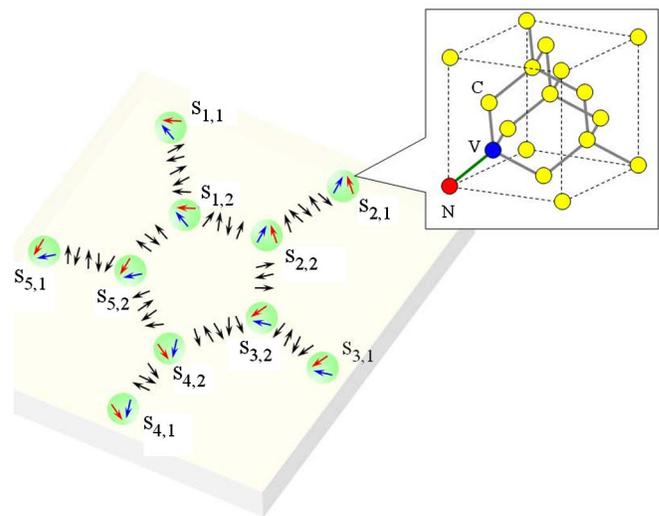


FIG. 1 (color online). A network of 10 NV centers (grouped in 5 chains) in a slab of diamond crystal (for illustration purposes only). Each NV center consists of a nuclear spin  $I = 1/2$  (red arrows, labeled as  $s_{i,j}$ ), serving for storage of quantum information, and an electronic spin  $S = 1$  (blue arrows), serving for transmission and readout of quantum information. The electronic spins in neighboring NV registers are coupled through the dark spin chain data bus (black arrows) formed by the dipole coupling of implanted nitrogen impurities.

[14]. Therefore, single-qubit and controlled-NOT (CNOT) operations, i.e., universal quantum gates [15] between the nuclear and electronic spins in a single NV center are accomplished. A related experiment was demonstrated in 2004 [16].

The electronic spins in neighboring NV centers are coupled through a dark spin chain data bus (DSCB). DSCB is composed of implanted nitrogen impurities coupled through dipole-dipole interactions. Very recently, Yao *et al.* showed that the states of the electronic spins in neighboring NV centers can be transferred through DSCB even if the states of DSCB are random [17]. Here we briefly review the mechanism of this scheme. The Hamiltonian for DSCB containing  $m$  nitrogen impurities can be effectively written in the form of an  $XX$  model:

$$H_{\text{DSCB}} = \sum_{i=1}^{m-1} 2\kappa(S_i^x S_{i+1}^x + S_i^y S_{i+1}^y), \quad (1)$$

where  $\kappa$  is the nearest-neighbor dipole-dipole coupling strength between impurity spins. Thus the total Hamiltonian describing the electronic spins in neighboring NV centers coupled through DSCB of  $m$  nitrogen impurities is

$$H = 2g(S_{\text{NV}_1}^x S_1^x + S_{\text{NV}_1}^y S_1^y + S_{\text{NV}_2}^x S_m^x + S_{\text{NV}_2}^y S_m^y) + H_{\text{DSCB}}, \quad (2)$$

where  $S_{\text{NV}_i}^{x,y}$  are the electronic spin operators of the NV center  $i$  ( $i = 1, 2$ ), and  $g$ , the coupling strength between the NV centers and DSCB, is controllable by utilizing the triplet energy-level structure of the electronic spins in NV centers [10].

The Hamiltonian  $H_{\text{DSCB}}$  can be transformed to a summation of  $m$  noninteracting fermionic modes with different energies [18], and the energy-level spacing of the electronic spins in NV centers can be tuned into resonance with a particular fermionic mode [17]. In this situation, the effective Hamiltonian governing the time evolution is

$$H_{\text{eff}} = t_k [f_k(c_0^\dagger + (-1)^{k-1} c_{m+1}^\dagger) + f_k^\dagger(c_0 + (-1)^{k-1} c_{m+1})], \quad (3)$$

where  $t_k = g\sqrt{\frac{2}{m+1}} \sin\frac{k\pi}{m+1}$ ,  $c_0$  and  $c_{m+1}$  are the fermionic operators associated with the electronic spins in neighboring NV centers:  $c_i = (S_i^x - iS_i^y) \prod_{j=0}^{i-1} (2S_j^z)$ , ( $i = 0, m+1$  correspond to  $\text{NV}_{1,2}$  respectively), and  $f_k^\dagger$  ( $f_k$ ) is the creation (annihilation) operator of the particular fermionic mode:  $f_k^\dagger = \sqrt{\frac{2}{m+1}} \sum_{i=1}^m \sin\frac{ik\pi}{m+1} c_i^\dagger$  ( $k$  is an integer and  $1 \leq k \leq m$ ). For a proper time evolution the states of the two boundary fermions are exchanged up to some phases ( $\pi$ ) depending on the states of all the fermions. The phases can be eliminated in a subsequent exchange of the states of boundary spins [17]. Furthermore, in a network of NV centers the direction of the state exchange (i.e., transfer) can be controlled by ensuring that the NV center in question is resonant only with the fermionic mode in the

desired direction, achieved with differing length of relevant DSCBs, i.e., different energy spectrums of fermionic modes. Recently, there have been experiments for studying the coupling between NV centers and the surrounding nitrogen impurities [19–21], which pave the way for realizing the state transfer between NV centers via DSCB.

Next we consider the realization of the HBT interferometer. First, we introduce a nonlocal Hadamard gate between the two nuclear spins of neighboring NV centers in the chain  $C_k$ , which is essential to the realization of the HBT interferometer.

$$U_{H,k} = \frac{1}{\sqrt{2}}(e^{i\phi_{10}}|10\rangle_k + e^{i\phi_{01}}|01\rangle_k)\langle 10|_k + \frac{1}{\sqrt{2}}(e^{i\phi'_{10}}|10\rangle_k + e^{i\phi'_{01}}|01\rangle_k)\langle 01|_k + |00\rangle_k\langle 00|_k + |11\rangle_k\langle 11|_k, \quad (4)$$

where  $|ij\rangle_k$  denotes the state of  $s_{k1}$  being  $|i\rangle$  and  $s_{k2}$  being  $|j\rangle$  [ $|0\rangle$  ( $|1\rangle$ ) is spin down (up)], and  $\phi'_{01} = \phi_{01} - \phi_{10} + \phi'_{10} + l\pi$  ( $l$  is an odd integer) by using  $U_{H,k}^\dagger U_{H,k} = \text{identity}$ . The gate  $U_{H,k}$  can be realized in the following way. We note that  $U_{H,k}$  is decomposed into several CNOT gates and single-spin operations [15]:  $U_{H,k} = XU_\alpha AYBYCX$ , where  $X = C_2\text{NOT}_1$  flips  $s_{k1}$  if  $s_{k2}$  is up,  $Y = C_1\text{NOT}_2$  is reversely defined,  $U_\alpha$  adds a phase  $e^{i\alpha}$  to  $s_{k1}$  if it is up [ $\alpha = (\phi_{01} + \phi'_{10} + \pi)/2$ ],  $A = R_z(\beta)R_y(\gamma/2)$ ,  $B = R_y(-\gamma/2)R_z(-(\delta + \beta)/2)$ ,  $C = R_z((\delta - \beta)/2)$  [ $R_i(\theta_0) = \exp(-i\theta_0\sigma_i/2)$ ,  $\sigma_i$  are Pauli matrices of  $s_{k2}$ ],  $\beta = \phi_{01} - \phi_{10}$ ,  $\gamma = \pi/2$ , and  $\delta = \phi'_{10} - \phi_{10} + \pi$ . Each operation in the decomposition can be performed; in particular, the CNOT gate is further decomposed into two local Hadamard gates with a controlled-Z gate in between [15]. The controlled-Z gate is realized by performing the same gate within a single NV center and transmitting states via DSCB [17].

The network is prepared in the initial state  $|10\rangle_1|10\rangle_2, \dots, |10\rangle_N$ , i.e., only the  $N$  outer nuclear spins are in the up state. We then perform  $U_{H,k}$  in Eq. (4) for all  $k$  so that the resulting state is

$$\prod_{k=1}^N \frac{e^{i\phi_{10}}|10\rangle_k + e^{i\phi_{01}}|01\rangle_k}{\sqrt{2}}. \quad (5)$$

In the network in the state (5), we rotate the states of the  $N$  inner nuclear spins counterclockwise through one inner-spin position, a process which we shall call counterclockwise rotation of inner nuclear spins or CRINS. This is achieved by sequentially exchanging the states of the  $N$  inner spins ( $N-1$ ) times:  $s_{1,2} \leftrightarrow s_{2,2}$ , followed by  $s_{2,2} \leftrightarrow s_{3,2}$  and so forth (see Fig. 1). The exchange of the states of two spins is similar to the process of realizing  $U_{H,k}$ , i.e., replace  $U_{H,k}$  with a swap gate (two CNOT gates). After this process is completed, the total state becomes

$$|\psi\rangle = \frac{1}{\sqrt{2^N}} \sum_{\{i_k\}} e^{i\xi_{i_1, i_2, \dots, i_N}} |i_1\rangle_{i_2}, \dots, |i_N\rangle, \quad (6)$$

where  $|i_k\rangle \in \{|00\rangle_k, |01\rangle_k, |10\rangle_k, |11\rangle_k\}$ . The values of  $\xi_{i_1, i_2, \dots, i_N}$  and  $i_k$  in (6) are not given explicitly as it is not

relevant to the present discussion. Next we consider a specific observable  $O_{r_k}^{(k)}$  defined for the  $k$ th spin chain ( $r_k$  indicates multiple choices)

$$O_{r_k}^{(k)} = \sum_{i,j} C_{r_k,i,j}^{(k)} |i\rangle_k \langle j|_k + \sum_{i',j'} C_{r_k,i',j'}^{(k)} |i'\rangle_k \langle j'|_k, \quad (7)$$

where  $i, j \in \{0, 01\}$  and  $i', j' \in \{00, 11\}$ . We note that  $O_{r_k}^{(k)}$  preserves the parity of the nuclear-spin excitations in the  $k$ th chain: two eigenvectors (say  $|\lambda_1^{(k)}\rangle, |\lambda_2^{(k)}\rangle$ ) of  $O_{r_k}^{(k)}$  are in the odd-parity subspace spanned by  $\{|10\rangle_k, |01\rangle_k\}$  and the other two (say  $|\lambda_3^{(k)}\rangle, |\lambda_4^{(k)}\rangle$ ) are in the even-parity subspace spanned by  $\{|00\rangle_k, |11\rangle_k\}$ , where  $\lambda_i^{(k)}$ 's are the eigenvalues of  $O_{r_k}^{(k)}$ . As  $O_{r_k}^{(k)}$  is Hermitian, i.e.,  $O_{r_k}^{(k)\dagger} = O_{r_k}^{(k)}$ , only the off-diagonal elements contain phases, e.g.,  $C_{r_k,01,10}^{(k)} = |C_{r_k,01,10}^{(k)}| e^{i\beta_{r_k}^{(k)}} = C_{r_k,10,01}^{(k)*}$  and  $C_{r_k,11,00}^{(k)} = |C_{r_k,11,00}^{(k)}| e^{i\beta_{r_k}^{(k)}} = C_{r_k,00,11}^{(k)*}$ . Let  $\mathcal{O} = \prod_{j=1}^{n_0} O_{r_{k_j}}^{(k_j)}$ , i.e., a joint operator of a subset of size  $n_0$  from the set  $\{O_{r_1}^{(1)}, O_{r_2}^{(2)}, \dots, O_{r_N}^{(N)}\}$ . Then it can be shown that the correlation function  $\langle \psi | \mathcal{O} | \psi \rangle$  can contain phases  $\beta_{r_k}^{(k)}$  and  $\beta_{r_k}^{(k)}$  only when  $n_0 = N$ . Furthermore,  $\langle \psi | \mathcal{O} | \psi \rangle$  can manifest itself as an intensity correlation function, i.e., the correlation function of the excitation numbers of the  $N$  outer spins in the network:  $\langle \psi' | \prod_{j=1}^{n_0} \sigma_{k_j,1}^\dagger \sigma_{k_j,1} | \psi' \rangle$ , where  $\sigma_{i,j} = (\sigma_{i,j}^x - i\sigma_{i,j}^y)/2$  and  $|\psi'\rangle = U_{r_k}^{(k)\dagger} |\psi\rangle$ . The unitary operator  $U_{r_k}^{(k)\dagger}$  consists of the four row eigenvectors of  $O_{r_k}^{(k)}$ :  $O_{r_k}^{(k)} = U_{r_k}^{(k)} \Lambda^{(k)} U_{r_k}^{(k)\dagger}$  with  $\Lambda^{(k)}$  a diagonal matrix in the basis  $|ij\rangle_k$ , ( $i, j \in \{0, 1\}$ ) encompassing the four eigenvalues of  $O_{r_k}^{(k)}$ . Moreover, the above two correlation functions also equal the following normal ordered  $n_0$ th-order intensity correlation function defined similarly as in quantum optics [22]:

$$G^{(n_0)}(k_{n_0}, \dots, k_1, k_1, \dots, k_{n_0}; t) = \langle \chi | \sigma_{k_{n_0},1}^\dagger(t), \dots, \sigma_{k_{1,1}}^\dagger(t) \sigma_{k_{1,1}}(t), \dots, \sigma_{k_{n_0},1}(t) | \chi \rangle, \quad (8)$$

where  $1 \leq k_1 < k_2 < \dots < k_{n_0} \leq N$ ,  $|\chi\rangle = \prod_{j=1}^{n_0} |10\rangle_{k_j}$  is the initial product state of the chains  $C_{k_1}, C_{k_2}, \dots, C_{k_{n_0}}$ ,  $\sigma_{k_j,1}(t)$  is the Heisenberg annihilation operator of the spin  $s_{k_j,1}$ , and  $t$  is the duration of all the previous operations performed on the system (including  $U_{H,k_j}$ , CRINS and  $U_{r_{k_j}}^{(k_j)\dagger}$ ). The measurement of the correlation function (8) can be realized through a repetitive fluorescence detection of the  $N$  outer spins [23,24]. In this perspective, we say that  $\langle \psi | \mathcal{O} | \psi \rangle$  fully manifests the HBT effect (see Supplemental Material [25] for an interpretation of the above discussion). An interesting case of such  $O_{r_k}^{(k)}$  is

$$O_{r_k}^{(k)} = \frac{1}{2} (e^{-i\beta_{r_k}^{(k)}} |10\rangle_k \langle 01|_k + e^{i\beta_{r_k}^{(k)}} |01\rangle_k \langle 10|_k + |10\rangle_k \langle 10|_k + |01\rangle_k \langle 01|_k + |11\rangle_k \langle 11|_k), \quad (9)$$

for which  $U_{r_k}^{(k)\dagger}$  can be calculated and is realized by performing  $U_{H,k}$  in Eq. (4) with  $\phi_{10} = \phi_{01} - \pi/2 = -\phi'_{10} + \pi/2 = \beta_{r_k}^{(k)}/2 + \pi/4$ .

The HBT effect originates from the interference of the different paths for the evolution of spin excitations. To see this, first we note that the Hadamard gate in Eq. (4) is analogous to the action of a beam splitter in optics [26], i.e., for  $|10\rangle_i \rightarrow (e^{i\phi_{10}} |10\rangle_i + e^{i\phi_{01}} |01\rangle_i)/\sqrt{2}$ , the initial excitation is equally reflected ( $e^{i\phi_{10}} |10\rangle_i$ ) and transmitted ( $e^{i\phi_{01}} |01\rangle_i$ ). Consider  $O_{r_k}^{(k)}$  to be Eq. (9). In this situation,  $\langle \psi' | \prod_{j=1}^{n_0} \sigma_{k_j,1}^\dagger \sigma_{k_j,1} | \psi' \rangle = \langle \psi | \prod_{j=1}^{n_0} O_{r_{k_j}}^{(k_j)} | \psi \rangle$ , as mentioned earlier. Since for  $n_0 < N$  this correlation function contains no phases  $\beta_{r_k}^{(k)}$  and  $\beta_{r_k}^{(k)}$ , we only consider the case of  $n_0 = N$ . It can be directly calculated and equals

$$\left[ \cos\left(\eta_0 + \sum_{k=1}^N \beta_{r_k}^{(k)}\right) + 1 \right] / 2^{2N-1}, \quad (10)$$

where  $\eta_0 = N(\phi_{10} - \phi_{01})$ . We note that only  $|10\rangle_1 |10\rangle_2, \dots, |10\rangle_N$  has a contribution. This state results from the interference of the following two paths: (1) the initial  $N$  outer spin excitations are reflected back twice by the ‘‘beam splitters’’ [see Fig. 2(a)]; and (2) they transmit through the ‘‘beam splitters’’ and the excitation in the chain  $C_k$  transfers to the chain  $C_{k-1}$  ( $C_0 \equiv C_N$ ) via CRINS. Subsequently, all the excitations transmit through the ‘‘beam splitters’’ and arrive at the outer spin positions [see Fig. 2(b)]. The two paths, possessing probability amplitudes  $2^{-N} \exp[i(N\phi_{10} + \sum_{k=1}^N \beta_{r_k}^{(k)}/2)] \equiv d_1$  and  $2^{-N} \exp[i(N\phi_{01} - \sum_{k=1}^N \beta_{r_k}^{(k)}/2)] \equiv d_2$ , respectively, interfere with each other because of the indistinguishability of the  $N$  excitations. The interference gives rise to  $|d_1 + d_2|^2$  for the probability of all the outer spins being up, which equals  $\langle \psi' | \prod_{k=1}^N \sigma_{k,1}^\dagger \sigma_{k,1} | \psi' \rangle$ .

The HBT effect can be considerably enhanced through a postselection process as shown below. Since  $O_{r_k}^{(k)}$  under the state  $|\psi\rangle$  can be measured through determining the four results of  $\{\sigma_{k,1}^\dagger \sigma_{k,1}, \sigma_{k,2}^\dagger \sigma_{k,2}\}$ , i.e.,  $\{i, j\}$  with  $i, j \in \{0, 1\}$  under  $|\psi'\rangle$ , one could locally determine whether the measurement results of the observable set  $\{O_{r_1}^{(1)}, O_{r_2}^{(2)}, \dots, O_{r_N}^{(N)}\}$  originate from the contributions of  $|\psi_G\rangle$  or  $|\psi_0\rangle$  (see Supplemental Material [25] for their expressions and the reason). Namely, if the contribution is from  $|\psi_0\rangle$ , there is at least one  $\{\sigma_{k,1}^\dagger \sigma_{k,1}, \sigma_{k,2}^\dagger \sigma_{k,2}\}$  for some  $k$  that has the result of either  $\{1, 1\}$  or  $\{0, 0\}$ . Therefore, by discarding these results and renormalizing the distribution of the retained results, one actually postselects the GHZ entanglement (i.e., the state  $|\psi_G\rangle$ ) [27]. For instance, choose  $O_{r_k}^{(k)}$  to be (9). Then  $\langle \psi | \prod_{j=1}^N O_{r_j}^{(j)} | \psi \rangle = \langle \psi' | \prod_{j=1}^N \sigma_{j,1}^\dagger \sigma_{j,1} | \psi' \rangle$  after postselection equals  $[\cos(\eta_0 + \sum_{k=1}^N \beta_{r_k}^{(k)}) + 1]/2^N$ , an exponential enhancement compared with (10). More generally, we note that a generic GHZ state is of the form  $(|\varphi_1\rangle|\varphi_2\rangle \dots |\varphi_N\rangle + |\varphi'_1\rangle|\varphi'_2\rangle \dots |\varphi'_N\rangle)/\sqrt{2}$ , where  $\langle \varphi_k | \varphi'_k \rangle = 0$  for all  $k$  [27]. Thus the state  $|\psi\rangle$  of Eq. (6) can be divided into  $2^{N-1}$  pairs of states and each pair is a GHZ-type state (see proof of the lemma in Supplemental Material [25]). One can postselect not only  $|\psi_G\rangle$  but also

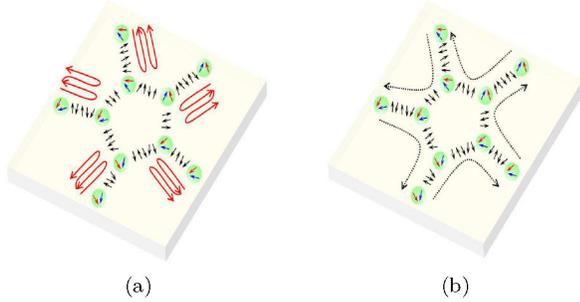


FIG. 2 (color online). The two paths that contribute to  $\langle \psi' | \otimes_{k=1}^N \sigma_{k,1}^\dagger \sigma_{k,1} | \psi' \rangle$ , ( $N = 5$ ).

any other types of GHZ states by appropriately extracting the measurement results of  $\{\sigma_{k,1}^\dagger \sigma_{k,1}, \sigma_{k,2}^\dagger \sigma_{k,2}\}$  for  $1 \leq k \leq N - 1$  (the  $N$ th result is determined by the previous ones). The correlation function (CF) for the state  $|\psi\rangle$  is a sum of the correlation functions (CF<sub>*i*</sub>) for  $2^{N-1}$  GHZ-type states with respective fractions  $P_i$  (the normalization constant of each postselected distribution):  $CF = \sum_{i=1}^{2^{N-1}} P_i CF_i$ . This is analogous to the decomposition of the interferometric fringe of the polychromatic light into those of monochromatic lights in a double-slit experiment.

The GHZ entanglement generated in the postselection process superposes two macroscopically distinct states as described in the previous paragraph. It shows the inconsistency of the elements-of-reality concept presented in the Einstein-Podolsky-Rosen paper [28] and has applications in multipartner quantum cryptography [29] and communication complexity tasks [30]. To test the nonlocality of GHZ entanglement, one could employ the Svetlichny inequality  $\langle S_N^\pm \rangle \leq 2^{N-1}$ , where  $S_N^\pm$  is generated via recursive relations  $S_N^\pm = S_{N-1}^\pm O_1^{(N)} \mp S_{N-1}^\mp O_2^{(N)}$  with  $O_{1,2}^{(N)}$  being two choices of  $O_{r_k}^{(N)}$  defined previously [31]. The  $N$ -particle Svetlichny inequality is maximally violated by the  $N$ -particle GHZ states and the violation rules out all hybrid local-nonlocal hidden variable models. Specifically, consider the expectation value of  $\prod_{k=1}^N O_{r_k}^{(k)}$ , i.e., the last equation of Sec. I of the Supplemental Material [25]. It can be verified that the equation is multiplied by  $2^{N-1}$  after postselection. Comparing it with Eq. (11) of Ref. [31], we find that the Svetlichny inequality can be maximally violated by a multiplication factor  $\sqrt{2}$  if we choose  $\beta_{r_k}^{(k)} = (\alpha_{r_k}^{(k)} - \eta_0/N)$ .

In conclusion, we have proposed a scheme to realize the multiparticle Hanbury Brown–Twiss interferometer in a network of NV centers. The interference effect can be enhanced considerably through the postselection to obtain multipartite Greenberger-Horne-Zeilinger entanglement whose nonlocality can be tested by Svetlichny inequality.

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