Orthogonally-Driven Superconducting Qubit in Circuit QED

M.J. Storcz*, 1 M. Mariantoni1, 2 H. Christ, 3 A. Emmert, 2 A. Marx, 2
W.D. Oliver, 4 R. Gross, 2 F.K. Wilhelm, 1 and E. Solano1, 5

1 Physics Department, CeNS and ASC, Ludwig-Maximilians-Universität, Theresienstrasse 37, 80333 Munich, Germany
2 Walther-Meissner-Institut, Bayerische Akademie der Wissenschaften, Walther-Meissner-Strasse 8, 85748 Garching, Germany
3 Max-Planck Institute for Quantum Optics, Hans-Kopfermann-Strasse 1, 85748 Garching, Germany
4 MIT Lincoln Laboratory, 244 Wood Street, Lexington, Massachusetts 02420, USA
5 Sección Física, Departamento de Ciencias, Pontificia Universidad Católica del Perú, Apartado 1761, Lima, Peru

(Dated: March 23, 2022)

We consider a superconducting charge qubit coupled to distinct orthogonal electromagnetic field modes belonging to a coplanar waveguide resonator and a microstrip transmission line. This architecture allows the simultaneous implementation of a Jaynes-Cummings and anti-Jaynes-Cummings dynamics, a resonant method for generating mesoscopic qubit-field superpositions and for field-state reconstruction. Furthermore, we utilize this setup to propose a field measurement technique that is, in principle, robust due to a fast pre-measurement to qubit dephasing and field relaxation.

PACS numbers: 74.81.Fa, 42.50.Dv, 32.80.-t

In recent years, superconducting quantum circuits have demonstrated key elements required for quantum information processing [1], including the possibility to prepare a desired qubit state, coherently manipulate it, read it out, and perform preliminary conditional gate operations [2, 3, 4]. In analogy to quantum-optical cavity QED, superconducting charge and flux qubits have been coupled to on-chip microwave resonators [5, 6], and universal two-qubit gates mediated by a single cavity mode have been proposed [7]. In such an architecture, the scaling of the system would require a homogeneous coupling of many qubits with the same cavity mode and a means to address each logical qubit individually. Prototypical advanced manipulation schemes have already been implemented in quantum-optical systems. In trapped ion experiments, tuned lasers can be switched between single-ion carrier excitations and ion-motion JC dynamics [8]. Also, in 3D microwave cavities, a flying atom can perform local rotations in Ramsey zones before and after entering the cavity, in which a Jaynes-Cummings (JC) interaction takes place [9]. In the emerging field of circuit QED [10, 11], a higher level of addressability and control is also desirable [12]. For instance, it is important to enable controlled intracavity qubit rotations, while keeping a switchable coupling to the cavity modes, aiming at intercavity qubit-qubit transfer of information. This implies the necessity to generate nonclassical field states and the possibility of measuring them via rapid qubit operations. These requirements could be simultaneously met if each qubit were coupled to two (or more) independent modes with orthogonal field polarizations.

In this Letter, we propose an architecture consisting of a superconducting charge qubit coupled to the quantized, discrete-mode spectrum of a quasi-1D coplanar waveguide (CWG) resonator, here called cavity [13], and to a multi-layer microstrip transmission line (MTL), which will be utilized to tune the qubit-cavity resonance. This archetypical construction ideally subjects the qubit to two orthogonal electromagnetic fields and constitutes the system for our theoretical investigations (See Fig. 1). We will show that strongly driving the MTL with coherent field pulses modulates the strength of the qubit-cavity interaction and, in the strong-driving limit [11], mediates the emergence of a simultaneous JC and anti-JC dynamics in the system, allowing the generation of nonclassical qubit-cavity states [12]. Furthermore, we propose the measurement of relevant observables of the microwave cavity field via a protocol that utilizes the MTL as an independent tool for read-out of the qubit population [13]. Finally, we will demonstrate that this technique is, in principle, robust due to a fast qubit-cavity pre-measurement to the presence of qubit dephasing and field decoherence processes.

The Hamiltonian describing the interaction between a single Cooper-pair box (CPB) charge qubit, the second harmonic of the undriven CWG resonator, and the MTL (see Fig. 1), driven with a propagating coherent state acting as an AC gate charge, can be written as [14]

\[
\hat{H} = -2E_C \left(1 - 2n_C^{DC} \right) \hat{\sigma}_z - E_1 \left(\frac{\Phi_x}{2} \right) \hat{\sigma}_x + \hbar \omega_C \hat{a}_2^\dagger \hat{a}_2
\]

\[+ \hbar g_{QC} \hat{\sigma}_z \left(\hat{a}_2^\dagger + \hat{a}_2\right) + 4E_C n_C^{AC} \left(t_s \right) \hat{\sigma}_z. \tag{1} \]

Here, \(E_C = e^2/C_S\) represents the charging energy of the CPB \((C_S\) is the total box capacitance), \(n_C^{DC}\) is the number of pairs induced by the DC gate voltage through the CWG cavity, \(E_1\) is the qubit Josephson energy, which can be tuned by an external quasi-static flux bias \(\Phi_x\), applied through a adequately engineered loop, \(\{\hat{a}_2^\dagger, \hat{a}_2\}\) are

*These authors contributed equally to this work.
the bosonic creation and destruction operators relative to the second harmonic of the CWG resonator, \( \omega_C \) and \( g_{QC} \) are the corresponding angular frequency and qubit-cavity vacuum Rabi coupling, respectively, \( n_{AC}^C(t) \) is the number of pairs induced by the AC gate voltage applied via the MTL, and \( \{ \hat{\sigma}_x, \hat{\sigma}_z \} \) are Pauli spin-1/2 operators. With \( 4E_Cn_{AC}^C(t) = (C_{QM}V_{QM}^0/2e)|\beta|\cos(\omega_M + \theta_{\beta}) \), we can define \( \hbar g_{QM}^\beta = (C_{QM}V_{QM}^0/2e)|\beta| \), where \( g_{QM} \) is the qubit-MLT coupling strength, \( C_{QM} \) and \( V_{QM}^0 \) are the associated capacitance and vacuum gate voltage, \( \beta = |\beta| \exp i\vartheta_{\beta} \) is the amplitude of the driven coherent field state \( |\beta\rangle \), \( \omega_M \) its angular frequency, and \( \vartheta_{\beta} \) a phase. The direct coupling between the CWG resonator and the MTL is suppressed due to the high-degree of isolation between their respective modes. The MTL magnetic field distribution is predominantly tangential to the CPB loop, suppressing unwanted jitters of the \( \hat{\sigma}_x \) term in Eq. (1) due to MTL noise.

Working in the eigenbasis \( \{|g\rangle, |e\rangle\} \) of the first two terms of Eq. (1), the system Hamiltonian takes the form

\[
\hat{H} = \frac{\Omega}{2} \hat{\sigma}_z + \hbar \omega_C \hat{a}_2^\dagger \hat{a}_2 + \left[ \hbar g_{QC} \left( \hat{a}_2^\dagger + \hat{a}_2 \right) + 4E_Cn_{AC}^C(t) \right] \times (\cos \theta \hat{\sigma}_z - \sin \theta \hat{\sigma}_x),
\]

where \( \{ \hat{\sigma}_x, \hat{\sigma}_z \} \) are Pauli matrices in the \( \{|g\rangle, |e\rangle\} \) eigenbasis, \( \Omega = \sqrt{E_j^2 + \left[ 4E_C (1-2n_{QC}^C) \right]^2} \) is the qubit level separation, and \( \theta = \arctan \left( E_j/4E_C (1-2n_{QC}^C) \right) \) is the mixing angle. Operating the qubit at the degeneracy point, i.e., for \( n_{QC}^C = 1/2 \), under complete resonance conditions \( |E_j/(2\hbar) = \omega_C = \omega_M| \), and within a standard rotating-wave approximation (RWA) in the interaction picture, Eq. (2) can be rewritten as

\[
\hat{H}^{\text{int}} = \hbar g_{QC} \left( \hat{\sigma}_+ \hat{a}_2^\dagger + \hat{\sigma}_- \hat{a}_2 \right) + \hbar g_{QM}^\beta \hat{\sigma}_x.
\]

A more advanced application of Eq. (3) is obtained in the strong-driving limit, \( g_{QM}^\beta \gg g_{QC} \). In this case, following Ref. [11], we utilize an additional interaction representation with respect to the second term on the r.h.s. of Eq. (3). Decomposing \( \sigma_\pm = (\sigma_x \pm i\sigma_y)/2 \) one allows to omit the quickly-precessing \( \sigma_y \) term through a second RWA, yielding the effective Hamiltonian

\[
\hat{H}_{\text{eff}} = \frac{\hbar}{2} g_{QC} \left( \hat{\sigma}_+ + \hat{\sigma}_- \right) \left( \hat{a}_2^\dagger + \hat{a}_2 \right).
\]

This strong-driving limit results in a circuit-QED realization of a simultaneously resonant JC and anti-JC superposition states between the qubit and the cavity field. For example, taking the initial qubit-cavity state to be \(|g, 0_C\rangle = (|+\rangle + |-\rangle)|0_C\rangle/\sqrt{2} \), where \(|\pm\rangle\) are the \( \hat{\sigma}_x \) eigenstates of the qubit with eigenvalues \( \pm 1 \), the evolution associated with Eq. (4) after an interaction time \( t_{\text{int}} \) yields the following Schrödinger cat state

\[
|\Psi_{\text{cat}}\rangle = \frac{(|+\rangle|\alpha\rangle + |-\rangle|-\alpha\rangle)}{\sqrt{2}},
\]

with \( \alpha = -i g_{QC} t_{\text{int}}/2 \). Using realistic experimental parameters, e.g., \( g_{QC} \approx 100 \text{ MHz} \) [18] and cavity decay rate \( \kappa_C \approx 0.1 \text{ MHz} \), would allow the generation of a Schrödinger cat state with amplitude \( |\alpha| = \sqrt{\langle \hat{n}_{\text{cat}} \rangle} \approx (g_{QC}/\kappa_C)^{1/3} \approx 8 \) (64 photons), obtained for a maximum interaction time \( t_{\text{int}} = 1/\kappa_C \), where \( \kappa_{\text{eff}} = |\alpha|^2 \kappa_C \). This amplitude compares well with the best values obtained in 3D microwave cavity QED with circular Rydberg atoms [19] and previous proposals in the resonant [20] and dispersive regimes [21].

In order to generate successfully the specific state \( |\Psi_{\text{cat}}\rangle \) inside the cavity, it is necessary to decouple the qubit-cavity system after the desired interaction time \( t_{\text{int}} \). This may be implemented by applying a second strong coherent state to the MTL, with decoupling frequency \( \omega_{\text{dec}} \) and amplitude \( |\gamma_{\text{dec}}| \), that AC-Stark shifts the CPB by an amount \( \delta \). At this point, a single-shot measurement of the qubit state \(|g\rangle \) \(|e\rangle\) would leave the field state in an even (odd) coherent state \( |\alpha\rangle = (|+\rangle + |-\rangle)/\sqrt{1 + \exp(-|\alpha|^2)} \langle \Psi_{\alpha}^+ \rangle = (|\alpha\rangle - |-\alpha\rangle)/\sqrt{1 - \exp(-|\alpha|^2)} \langle \Psi_{\alpha}^- \rangle \). This qubit measurement may be made, for instance, by driving a probing field via the third harmonic of the CWG resonator, which is strongly detuned from the qubit transition frequency by an amount \( \omega_{\text{C}}/2 \gg g_{QC} \), for a time \( t_{\text{meas}} \), and thereby allowing an independent QND read-out of
the qubit population via the resultant phase shift of the probing field \(22\).

The architecture proposed here could be applied directly to implement conventional cavity field measurements, consisting of an initial qubit-cavity pre-measurement in which the qubit acts as the quantum probe with which the cavity field is entangled, followed by a measurement of the qubit \(23\). Typically, one requires an interaction time sufficiently long (order of the inverse interaction frequency) to entangle strongly the cavity field with the qubit. Accordingly, one typically desires to operate in the qubit-cavity strong-coupling regime to minimize this interaction time and, thereby, reduce the noisy action of decoherence \(8,9\). In contrast to this common approach, we will show below an alternative means to implement a measurement of the cavity field with a relatively fast qubit-cavity pre-measurement (small fraction of the inverse interaction time) and minimal action of decoherence \(13\).

We consider now the even coherent state \(|\Psi_0^e\rangle\) obtained after measuring the qubit in the ground state \(|g\rangle\), as seen from Eq. \(5\). Through the MTL, a \(\pi/2\)-pulse of duration \(t_{\text{rot}}\) and resonant with the shifted qubit transition can be applied, yielding \(\rho_0(t_S) = |+\rangle\langle +| \otimes \rho_0^e\), with \(|+\rangle = (|g\rangle + e^{i\phi}|e\rangle)/\sqrt{2}\). Here, \(\phi\) is a relative phase, \(\rho_0^e = |\Psi_0^e\rangle\langle \Psi_0^e|\), and \(t_S = t_{\text{int}} + t_m + t_{\text{rot}}\) may be set to \(t_S = 0\). At this point, if the aforementioned decoupling pulse (frequency \(\omega_{\text{dec}}\) and amplitude \(|\gamma_{\text{dec}}|\)) applied via the MTL is turned off for a time \(t_{\text{off}}\), the qubit-cavity state \(\hat{\rho}_S(0)\) will evolve via a resonant JC dynamics for a dimensionless time \(\tau = g_{QC}t_{\text{off}}\).

In the presence of a dispersive bath, producing qubit dephasing, and a thermal bath at zero temperature, inducing field dissipation, the system dynamical equation becomes

\[
\frac{d}{d\tau} \rho_0 = \frac{1}{i\hbar} \left[ \hat{H}_{QC}, \rho_0 \right] - \sum_i \frac{\gamma_i}{2} \left( \left\{ \hat{A}_i, \hat{A}_i^\dagger \rho_0 \right\} - 2 \hat{A}_i \rho_0 \hat{A}_i^\dagger \right),
\]

where \(\hat{H}_{QC} = \hbar g_{QC} \left( \hat{\sigma}^+ \hat{a}_2 + \hat{\sigma}^- \hat{a}_2^\dagger \right)\), the braces denote anti-commutators, and the sum runs over two indices corresponding to qubit dephasing \(\langle \hat{A}_1 = \hat{\sigma}^z, \gamma_1 = \gamma_\phi \rangle\) and cavity decay \(\langle \hat{A}_2 = \hat{a}_2, \gamma_2 = \kappa_c \rangle\), respectively. Above, we have neglected the energy relaxation of the qubit which is assumed to happen on a longer timescale \(22\). Calculating now \(dP_e(\tau)/d\tau = d\langle |e\rangle \langle e| \rangle/d\tau = \text{Tr} \left( \dot{\rho}_S |e\rangle \langle e| \right)\), and after some algebra, we obtain

\[
\frac{d}{d\tau} P_e(\tau) = \frac{1}{i\hbar} \left( \langle |e\rangle \langle e|, \hat{H}_{QC} \right). \quad (6)
\]

This expression does not contain terms involving qubit dephasing or field decay rates, which, ideally, are eliminated by the expectation value calculated at time \(\tau\). Evaluating Eq. \(6\) in the limit \(\tau \to 0^+\) yields the useful result

\[
\langle \dot{Y}_\phi \rangle = \left. \frac{d}{d\tau} P_e(\tau) \right|_{\tau \to 0^+}, \quad (7)
\]

where \(\dot{Y}_\phi = \left( \hat{a}_2 e^{-i\phi} - \hat{a}_2^\dagger e^{i\phi} \right)/2i\) is a field quadrature, conjugate to the quadrature \(\dot{X}_\phi = \left( \hat{a}_2 e^{-i\phi} + \hat{a}_2^\dagger e^{i\phi} \right)/2\) obtained by replacing \(\phi \to (\phi - \pi/2)\). Equation \(7\) shows mathematically that the first derivative of the measured excited-state qubit population obtained at infinitesimally small interaction time contains information about cavity-field observables with no influence of decoherence processes. If the qubit is now prepared in the excited state \(|e\rangle\), it is also possible to determine the mean photon number of the cavity field, via

\[
\langle \hat{n} \rangle = \langle \hat{a}^\dagger \hat{a} \rangle = \frac{1}{2} \frac{d^2}{d\tau^2} P_e(\tau) \bigg|_{\tau \to 0^+} - 1. \quad (8)
\]

Fast pre-measurements, plus the near unit visibility read-out of qubit populations reported in Ref. \(22\), turns the proposed measurement scheme into an alternative means to measure cavity-field observables with reduced noise disturbance.

Although Eqs. \(7\) and \(8\) are exact mathematical expressions, they do not represent a realistic theoretical description of a measurement. In this sense, it is useful to consider an estimator for the derivatives in Eqs. \(7\) and \(8\) over short, but non-zero, measurement times \(\Delta\tau\) through the Taylor expansion: \(P(\tau + \Delta\tau) = P(\tau) + P'(\tau)\Delta\tau + P''(\tau)(\Delta\tau)^2/2! + O((\Delta\tau)^3)\). Subsequently, it can be shown that

\[
\frac{\Delta P_e^{+e}(\tau)}{\Delta\tau}
\]

\(\approx \langle \dot{Y}_\phi \rangle - \left[ 1 + \frac{\kappa_c + \gamma_\phi}{4g_{QC}} \langle \dot{Y}_\phi \rangle \right] \Delta\tau, \quad (9)
\]

where the second term on the r.h.s. is the dominant higher-order contribution to Eq. \(7\), which, even for ideal ensemble averaging, contains the field decay rate \(\kappa_c\) and the qubit dephasing rate \(\gamma_\phi\). Equation \(9\) shows that in the strong-coupling regime, \(\{\kappa_c, \gamma_\phi\} < g_{QC}\), or even in the weak-coupling regime, \(\{\kappa_c, \gamma_\phi\} > g_{QC}\), it is possible to identify a time \(\Delta\tau\) that is long enough to allow the readout of \(\langle \dot{Y}_\phi \rangle\) and short enough to suppress the effects of decoherence during the pre-measurement. This result illustrates that, in principle, strong-coupling between the probe and the system is not required to implement the proposed measurement scheme.

Establishing an accurate estimator with the prescribed accuracy for the ensemble averages in Eqs. \(7\) and \(8\) will, in general, require many repetitions of the prescribed measurement. There is a trade-off between the length of the physically implemented \(\Delta\tau\), the number of measurement repetitions, and the strength of the qubit-cavity coupling required to achieve a desired degree of...
noise immunity. In consequence, the proposed measurement technique consists of two clear steps. First, a fast pre-measurement allows the pointer (qubit) to encode the information about the system (cavity field), minimizing decoherence processes. Second, the readout of the pointer, happening typically over longer times, completes the quantum measurement procedure.

Finally, another useful application of our scheme is the possibility to perform full state reconstruction of an unknown field state \( \hat{\rho} = |\Psi\rangle \langle \Psi | \) or, equivalently, of its corresponding characteristic function \( \chi(\hat{\alpha}) = \text{Tr} [\hat{\rho} D(\hat{\alpha})] \). Here, \( D(\hat{\alpha}) \) is the displacement operator and \( \hat{\alpha} \) the complex amplitude in phase space of an arbitrary coherent state \([24]\). For the sake of convenience, a pure field state will be considered here, even though these results remain valid for any arbitrary mixed state. Following a similar protocol to the one outlined above, the qubit can be initially prepared in the state \( |+\theta\rangle = (|+\rangle + e^{i\theta} |\rangle)/\sqrt{2} \), while the CWG cavity, populated with the unknown field \( \langle \Psi_C \rangle \), stays unperturbed. At this point, the interaction described in Eq. (4) can be turned on and the initial state \( |+\theta\rangle |\Psi_C \rangle \), after an interaction time \( t_d \), evolves to

\[
|\Psi(\hat{\alpha}, \theta)\rangle = (|+\rangle D(\hat{\alpha}) |\Psi_C \rangle + e^{i\theta} |\rangle D(-\hat{\alpha}) |\Psi_C \rangle) / \sqrt{2},
\]

with \( \hat{\alpha} = -i g Q t_d /2 \). By measuring the ground state qubit population \( P_g(\hat{\alpha}, \theta) \), given that the initial state was \( |+\theta\rangle \), we can retrieve the characteristic function through

\[
\chi(\alpha) = \left[ P_g \left( \frac{\alpha}{2}, 0 \right) - \frac{1}{2} \right] + i \left[ P_g \left( \frac{\alpha}{2}, \frac{\pi}{2} \right) - \frac{1}{2} \right].
\]

From this measured function \( \chi(\alpha) \), it is possible to derive \( \hat{\rho} \), or its associated Wigner function \([21][22]\), via a Fourier transform \([24]\), achieving a full-state reconstruction.

One means of coupling a CPB to two quasi-orthogonal electric fields is to make use of a multi-layer technology. In Fig. 1 dielectric Layer 0 is the substrate for the entire structure, and the MTL ground plane is Metal 0. Dielectric Layer 1 serves as a substrate for the CWG resonator-CPB structures, which are made from Metal 1. Dielectric Layer 2 supports the MTL made from Metal 2. The proper engineering of this structure results in quasi-perpendicular electric fields at the CPB. These fields are described by dyadic Green’s functions and can be evaluated numerically with the method of moments \([26]\). The results of our simulations indicate about \(-40\) dB isolation at 5 GHz. For example, driving the MTL with a coherent state of \( |\hat{a}\rangle \approx 10^3 \) photons would populate the CWG resonator with \( \sim 0.1 \) photons.

In conclusion, we have considered a CPB coupled to the orthogonal modes of a CWG and MTL resonators. This architecture allows the engineering of simultaneous JC and anti-JC dynamics, which we utilized to generate mesoscopic entangled states of the CPB and the CWG. It may be applied to the measurement of the field quadratures using a measurement scheme requiring a relatively short interaction time between the system (cavity) and the probe (qubit). We have also shown that full state reconstruction of unknown fields can be obtained. The multimode concepts described here should bring greater flexibility to the field of circuit QED, in particular those aiming at the systematic scaling-up of quantum gates and quantum information devices.

This work was supported by DFG through SFB 631. The work at Lincoln Laboratory was sponsored by the US DoD under Air Force Contract No. FA8721-05-C-0002. ES acknowledges support of EU EuroSQIP project.

[1] M.A. Nielsen and I.L. Chuang, Quantum Computation and Quantum Information, (Cambridge University Press, Cambridge, 2000).
[2] Y. Makhlin, G. Schön, and A. Shnirman, Rev. Mod. Phys. 73, 357 (2001).
[3] M.H. Devoret, A. Wallraff, and J.M. Martinis, cond-mat/0411174 (2004).
[4] T. Yamamoto et al., Nature 425 941 (2003).
[5] A. Wallraff et al., Nature 431, 162 (2004).
[6] J. Johansson et al., Phys. Rev. Lett. 96, 127006 (2006).
[7] A. Blais et al., Phys. Rev. A 69, 062320 (2004).
[8] D. Leibfried, R. Blatt, C. Monroe, and D. Wineland, Rev. Mod. Phys. 75, 281 (2003).
[9] J.M. Raimond, M. Brune, and S. Haroche, Rev. Mod. Phys. 73, 565 (2001).
[10] G. Johansson, L. Tornberg, and C. M. Wilson, Phys. Rev. B 74, 100504(R) (2006).
[11] E. Solano, G.S. Agarwal, and H. Walther, Phys. Rev. Lett. 90, 027903 (2003).
[12] M. Brune et al., Phys. Rev. Lett. 77, 4887 (1996).
[13] P. Lougovski, H. Walther, and E. Solano, Eur. Phys. J. D 38, 423 (2006).
[14] Constant energy offsets, the cavity vacuum term, and negligible second order terms have been omitted.
[15] T. Itoh, Numerical Techniques for Microwave and Millimeter-Wave Passive Structures, John Wiley & Sons, Inc. (New York), 1989.
[16] M. Mariani et al., arXiv:cond-mat/0509757 (2005).
[17] D.I. Schuster et al., Phys. Rev. Lett. 94, 123602 (2005).
[18] D.I. Schuster, private communication.
[19] A. Auffeves et al., Phys. Rev. Lett. 91, 230405 (2003).
[20] Y.-x. Liu, L. F. Wei, and F. Nori, Phys. Rev. A 71, 063820 (2005).
[21] F. de Melo, L. Aolita, F. Toscano, and L. Davidovich, Phys. Rev. A 73, 030303 (2006).
[22] A. Wallraff et al., Phys. Rev. Lett. 95, 060501 (2005).
[23] W.H. Zurek, Rev. Mod. Phys. 75, 715 (2003).
[24] K.E. Cahill and R.J. Glauber, Phys. Rev. 177, 1857 (1969).
[25] X. Zou, K. Pahlke, and W. Mathis, Phys. Rev. A 69, 015802 (2004).
[26] R.E. Collin, Field Theory and Guided Waves, Wiley-IEEE Press, 2nd Edition (New York), 1990.