ac Josephson effect in the resonant tunneling through mesoscopic superconducting junctions (postprint)

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Full Text

Preamble

ac Josephson effect in the resonant tunneling through mesoscopic superconducting junctions

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Abstract

We investigate the ac Josephson effect in resonant tunneling through mesoscopic superconducting junctions. In the presence of microwave irradiation, we show that the trajectory of multiple Andreev reflections can be closed by emitting or absorbing photons. Consequently, photon-assisted Andreev states are formed and play the role of carrying supercurrent. On the Shapiro steps, a dc component appears when the resonant level E_0 is near the position $V_L+\frac{1}{2}m_1\omega = V_R+\frac{1}{2}m_2\omega$ ($e=\hbar=1$), where m_1 and m_2 are integers, ω is the frequency of the microwave, and V_L and V_R are the chemical potentials of the superconducting leads. An analytical result is derived in the limit $\Delta \to \infty$, based on which new features of the ac Josephson effect are revealed.

Introduction

Since the achievement of atom-size superconducting quantum point contact (SQPC) by break junction technique [?, ?, ?, ?], considerable theoretical work has been devoted to the transport problem in S-I-S or S-N-S structures (S=superconductor, I=tunnel barrier, N=normal metal). Comparison between experimental data and theoretical calculation suggests that the coherent picture of multiple Andreev reflection (MAR) is the central concept in understanding transport through such mesoscopic SQPCs. Nevertheless, less attention has been paid to resonant tunneling (especially in ac transport) through mesoscopic superconducting junctions, in which the central region consists of discrete electronic states. Such resonant tunneling can be achieved in nanoparticles sandwiched between superconducting membranes [?], metallic clusters absorbed on STM substrates [?], a piece of carbon nanotube tunnel-coupled to superconducting electrodes [?], or quantum dots fabricated in S/2DEG hybrid structures [?], among others.

The transport through these structures is greatly modified due to the existence of discrete energy levels in the central region. It has been shown by Yeyati et al. [?] and Johansson et al. [?] that the combination of MAR processes and resonant transmission gives rise to a rich subharmonic gap structure (SGS). Consider the transport through a resonant level with width Γ coupled to superconducting leads with gap Δ. In the limit $\Gamma \gg \Delta$, the resonance is sufficiently broadened such that the I-V characteristics mimic those of SQPC. In the regime $\Gamma \sim \Delta$, pronounced SGS appears with concomitant negative differential conductance, which distinguishes it from nonresonant transport through SQPC. In the limit $\Gamma \ll \Delta$, although SGS is more particular in a log-scale plot, the tunnel current as a whole is exponentially small. This is because a narrow resonance cannot cover the entire MAR trajectory, and it is very unlikely that multiple discrete levels are positioned exactly where the MAR trajectory passes (see Fig. 1a).

In this paper, we address the effect of microwave (MW) irradiation on resonant tunneling through mesoscopic superconducting junctions. We show that the tunnel current is greatly enhanced on the Shapiro steps when the resonant level E_0 is near the position $V_L+\frac{1}{2}m_1\omega = V_R+\frac{1}{2}m_2\omega$ ($e=\hbar=1$ throughout), where m_1 and m_2 are integers, ω is the frequency of the MW field, and V_L and V_R are the chemical potentials of the left and right superconducting leads. This can be attributed to the formation of photon-assisted Andreev states (PAAS), which play the role of carrying supercurrent. In the limit $\Delta \to \infty$, analytical results are derived, revealing new features of the ac Josephson effect in the case of resonant tunneling.

Hamiltonian and Formulation

The model Hamiltonian reads $H = \sum_{\beta=L,R} H_\beta + H_{\text{cent}} + H_T$, in which $H_\beta = \sum_{k\sigma} \varepsilon_k c^\dagger_{\beta k\sigma} c_{\beta k\sigma} + (\Delta e^{i\phi_\beta} \sum_k c^\dagger_{\beta -k\downarrow} c^\dagger_{\beta k\uparrow} + \text{H.c.})$ is the Hamiltonian for the βth superconducting lead, $H_{\text{cent}} = \sum_\sigma E_0 a^\dagger_\sigma a_\sigma$ describes the resonant level in the central region, and $H_T = \sum_{k\sigma} (V_\beta + W_\beta \cos\omega t)(c^\dagger_{\beta k\sigma} a_\sigma + \text{H.c.})$ represents the tunnel couplings. Here $\tilde{V}_\beta(t) \equiv V_\beta + W_\beta \cos\omega t$ is the time-dependent voltage drop across the βth barrier, where V_β is the chemical potential controlled by dc bias voltage and $W_\beta \cos\omega t$ is the ac voltage induced by MW irradiation. For simplicity, we assume that Coulomb interaction and multi-level effects can be ignored in the central region. This assumption is somewhat too ideal to achieve experimentally, but allows us to obtain analytical results and is therefore instructive for understanding more complicated cases.

Keldysh Green' s functions are defined in the 2×2 Nambu representation:

$$G^{r,a,<}_{t_1 t_2} \equiv \begin{pmatrix} \langle\langle a_\uparrow(t_1) | a^\dagger_\uparrow(t_2) \rangle\rangle^{r,a,<} & \langle\langle a_\uparrow(t_1) | a_\downarrow(t_2) \rangle\rangle^{r,a,<} \\ \langle\langle a^\dagger_\downarrow(t_1) | a^\dagger_\uparrow(t_2) \rangle\rangle^{r,a,<} & \langle\langle a^\dagger_\downarrow(t_1) | a_\downarrow(t_2) \rangle\rangle^{r,a,<} \end{pmatrix}$$

The time-dependent current flowing out of the βth lead can be expressed as

$$I_\beta(t) = 2e\, \text{Re}\, \text{Tr}\, \sigma_z \left([G^r \cdot \Sigma^<_\beta]_{tt} + [G^< \cdot \Sigma^a_\beta]_{tt} \right)$$

in which σ_z is the third Pauli matrix, the shorthand notation $[\cdots]_{tt}$ stands for integration over the intermediate time variable, and Σ_β is the self-energy due to tunnel coupling between the central region and the βth lead. G^r, G^a, and $G^<$ satisfy the integral equations:

$$G^{r,a} = g^{r,a} + g^{r,a} \cdot \Sigma^{r,a} \cdot G^{r,a}, \quad G^< = G^r \cdot \Sigma^< \cdot G^a$$

where $g^{r,a}_{t_1 t_2} = e^{-i\int_{t_2}^{t_1} \tilde{V}_\beta(t')dt'} g^{r,a}(\epsilon) e^{i\int_{t_2}^{t_1} \tilde{V}_\beta(t')dt'}$ and $\Sigma^{r,a,<}_{\beta,t_1 t_2} = U_\beta(t_1) e^{-i\epsilon(t_1-t_2)} \Sigma^{r,a,<}_\beta(\epsilon) U^\dagger_\beta(t_2)$, in which

$$g^{r,a}(\epsilon) = \begin{pmatrix} \frac{1}{\epsilon - E_0 \pm i0^+} & 0 \\ 0 & \frac{1}{\epsilon + E_0 \pm i0^+} \end{pmatrix}$$

$$\Sigma^{r,a}_\beta(\epsilon) = \Gamma_\beta \begin{pmatrix} \frac{\epsilon}{\sqrt{\epsilon^2 - \Delta^2}} \pm i\eta & -\frac{\Delta e^{-i\phi_\beta}}{\sqrt{\epsilon^2 - \Delta^2}} \\ -\frac{\Delta e^{i\phi_\beta}}{\sqrt{\epsilon^2 - \Delta^2}} & \frac{\epsilon}{\sqrt{\epsilon^2 - \Delta^2}} \pm i\eta \end{pmatrix}$$

$$\Sigma^<_\beta(\epsilon) = f(\epsilon)(\Sigma^r_\beta(\epsilon) - \Sigma^a_\beta(\epsilon))$$

with Γ_β being the coupling strength, η the dephasing rate in the superconducting lead, $f(\epsilon)$ the Fermi distribution function, and $\mathrm{Im}\sqrt{z} > 0$ as a convention. The remaining task is to solve these integral equations and evaluate the dc component of the time-dependent current.

There are two intrinsic frequencies in the problem, $\omega_1 = 2eV = 2(V_L - V_R)$ and $\omega_2 = \omega$. Generally, one may perform a Fourier transform $A_{t_1 t_2} = e^{i(l_1\omega_1 + l_2\omega_2)t_1} \sum_{l_1 l_2} \int \frac{d\epsilon}{2\pi} e^{-i\epsilon(t_1 - t_2)} \tilde{A}_{l_1 l_2}(\epsilon)$ and derive recursive relations for $\tilde{A}_{l_1 l_2}(\epsilon)$, as done in [?]. The calculation in this way relies on numerical computing power, and analytical results are impossible. We note, however, that in the case of narrow resonance, the dc component appears only if $\omega_1 = N\omega_2$ with N being an integer. This becomes clear by considering PAAS shown in Figs. 1b and 1c: an electron (or hole) enters through the resonant level E_0 and is Andreev reflected by the right superconducting gap; the reflected hole (or electron) exchanges m_1 photons with the MW field so that it can again pass through E_0. Then the hole (or electron) is Andreev reflected by the left superconducting gap as an electron (or hole), and exchanges m_2 photons to close the trajectory (detailed discussion of photon-assisted Andreev reflection is available in [?]). It is easy to see that the formation of PAAS requires $2V = (m_1 + m_2)\omega$. For this reason, we shall only consider the case of $V = V_N = N\omega/2$, while the current deviated from this condition is negligibly small.

The problem is largely simplified since ω can be used as the basic frequency in the Fourier expansion. Define the Fourier transformation as

$$A_{t_1 t_2} = \sum_l e^{il\omega t_1} \int \frac{d\epsilon}{2\pi} e^{-i\epsilon(t_1 - t_2)} \tilde{A}_l(\epsilon), \quad A_{mn}(\epsilon) = \tilde{A}_{m-n}(\epsilon + m\omega)$$

The definition guarantees the nice property that if $C = A \cdot B$, then $C_{mn}(\epsilon) = \sum_k A_{mk}(\epsilon) B_{kn}(\epsilon)$. The Fourier-transformed $g^{r,a}$ and $\Sigma^{r,a,<}$ are:

$$g^{r,a}_{mn}(\epsilon) = \delta_{mn} g^{r,a}(\epsilon + m\omega)$$

$$\Sigma^{r,a,<}_{R,mn}(\epsilon) = J_{l-m}(\alpha_R)\Sigma^{r,a,<}_{R,11}(\epsilon^0_l)J_{l-n}(\alpha_R) + J_{m-l}(\alpha_R)\Sigma^{r,a,<}_{R,21}(\epsilon^0_l)J_{l-n}(\alpha_R) + J_{l-m}(\alpha_R)\Sigma^{r,a,<}_{R,12}(\epsilon^0_l)J_{n-l}(\alpha_R) + J_{m-l}(\alpha$$

$$\Sigma^{r,a,<}_{L,mn}(\epsilon) = J_{l-m}(\alpha_L)\Sigma^{r,a,<}_{L,11}(\epsilon^+_l)J_{l-n}(\alpha_L) + J_{m-l-N}(\alpha_L)\Sigma^{r,a,<}_{L,21}(\epsilon^-_l)J_{l-n}(\alpha_L) + J_{l-m}(\alpha_L)\Sigma^{r,a,<}_{L,12}(\epsilon^+_l)J_{n-l}(\alpha_L) + J_{m-}$$

in which $\alpha_\beta \equiv W_\beta/\omega$ is the MW strength on the βth tunnel barrier, $J_n(x)$ the nth Bessel function, $\epsilon_l^0 = \epsilon + l\omega$, $\epsilon_l^\pm = \epsilon + (l \pm \frac{1}{2}N)\omega$, $V_N = N\omega/2$, $V_L = V_N$ and $V_R = 0$ are set as a convention.

Correspondingly, the equations for G^r, G^a, and $G^<$ are Fourier transformed into

$$G^{r,a}(\epsilon) = g^{r,a}(\epsilon) + g^{r,a}(\epsilon)\Sigma^{r,a}(\epsilon)G^{r,a}(\epsilon)$$

and

$$G^<(\epsilon) = G^r(\epsilon)\Sigma^<(\epsilon)G^a(\epsilon)$$

or equivalently,

$$G^{r,a}(\epsilon) = g^{r,a}(\epsilon) + g^{r,a}(\epsilon)\Sigma^{r,a}(\epsilon)g^{r,a}(\epsilon) + g^{r,a}(\epsilon)\Sigma^{r,a}(\epsilon)g^{r,a}(\epsilon)\Sigma^{r,a}(\epsilon)g^{r,a}(\epsilon) + \cdots$$

and

$$G^<(\epsilon) = [g^r(\epsilon) + g^r(\epsilon)\Sigma^r(\epsilon)g^r(\epsilon) + \cdots]\Sigma^<(\epsilon)[g^a(\epsilon) + g^a(\epsilon)\Sigma^a(\epsilon)g^a(\epsilon) + \cdots]$$

We note that finite-order perturbation expansion is inadequate for this problem because the formation of PAAS involves up to infinite order of tunneling processes. To resum the series, we adopt the resonant approximation [**?**]:

$$\frac{1}{\epsilon + l_1\omega - E_0 + i0^+} \cdot \frac{1}{\epsilon + l_2\omega - E_0 + i0^+} \approx \delta_{l_1 l_2} \frac{1}{(\epsilon + l_1\omega - E_0 + i0^+)^2}$$

The approximation implies that the overlap between sidebands $E_0 + l_1\omega$ and $E_0 + l_2\omega$ can be ignored if $l_1 \neq l_2$, which is justified when $\Gamma_\beta \ll \omega$. Applying the approximation to the equations above, one can obtain the solution

$$G_{mn}^{r,a,<}(\epsilon) = \begin{pmatrix} G_{mn,11}^{r,a,<} & G_{mn,12}^{r,a,<} \\ G_{mn,21}^{r,a,<} & G_{mn,22}^{r,a,<} \end{pmatrix}$$

in which

$$G_{mn}^{r,a,<}(\epsilon) = \frac{1}{(\epsilon + m\omega - E_0 \pm i0^+)^2} D_{mn}^{r,a,<}(\epsilon)$$

with

$$D^{r,a}_{mn}(\epsilon) = \begin{pmatrix} \epsilon + m\omega - \Sigma^{r,a}_{mm,11}(\epsilon) & -\Sigma^{r,a}_{mn,12}(\epsilon) \\ -\Sigma^{r,a}_{nm,21}(\epsilon) & \epsilon + n\omega - \Sigma^{r,a}_{nn,22}(\epsilon) \end{pmatrix}^{-1}$$

and

$$D^{<}_{mn}(\epsilon) = \begin{pmatrix} \Sigma^{<}_{mm,11}(\epsilon) & \Sigma^{<}_{mn,12}(\epsilon) \\ \Sigma^{<}_{nm,21}(\epsilon) & \Sigma^{<}_{nn,22}(\epsilon) \end{pmatrix}$$

It can be shown that the relations $G^r = (G^a)^\dagger$ and $G^< = -(G^<)^\dagger$ still hold in the solution.

With these Green' s functions, the time-dependent current can be expressed as a summation over ac components:

$$I_\beta(t) = 2e\,\mathrm{Re}\sum_l e^{il\omega t}\int \frac{d\epsilon}{2\pi}\,\mathrm{Tr}\,\sigma_z\left[G^r(\epsilon)\Sigma^<_\beta(\epsilon) + G^<(\epsilon)\Sigma^a_\beta(\epsilon)\right]_{l0}$$

The current formula can be applied to the study of ac harmonics. However, we are more interested in the dc component $\bar{I} = \bar{I}_L = -\bar{I}_R$, due to experimental considerations. To produce analytical results, we take the limit $\Delta \to \infty$, for which all single-particle processes are forbidden and Andreev reflection is the only conducting mechanism. After some algebra and taking the limit $\eta \to 0$ (noting that $\lim_{\eta\to 0}\frac{\eta}{x^2+\eta^2} = \pi\delta(x)$), one can derive

$$\bar{I} = 2e\frac{\Delta^2}{4}\sin\phi\sum_m J_{m-N}(2\alpha_L)J_m(2\alpha_R)\int\frac{d\epsilon}{2\pi}f(\epsilon)\left[\frac{1}{(\epsilon - E^+_m)^2 + |\Gamma_m|^2} - \frac{1}{(\epsilon - E^-_m)^2 + |\Gamma_m|^2}\right]$$

where

$$E^\pm_m = E_0 - \frac{V_N}{2} \pm \sqrt{\left(\frac{V_N}{2}\right)^2 + |\Gamma_m|^2}$$

and

$$\Gamma_m = \frac{\Delta}{2}\left[\Gamma_R J_m(2\alpha_R) + \Gamma_L e^{-i\phi}J_{m-N}(2\alpha_L)\right]$$

with $\phi = \phi_L - \phi_R$ being the phase difference between the two superconductors.

Equation (26) is the central result of this paper, which gives the dc component of the Josephson current in resonant tunneling through mesoscopic superconducting junctions.

Before presenting numerical studies, we make a few remarks on this result: (1) The phase dependence of $\bar{I}$ is mainly determined by the prefactor $\sin\phi$, and a weak $\cos\phi$ dependence is hidden in Γ_m. For this reason, ϕ is set as $\pi/2$ in the numerical calculation. (2) For the special case $N=0$, $\alpha_L=\alpha_R=0$, and $\Gamma_L=\Gamma_R=\Gamma$, $\bar{I}=2e\frac{\Delta^2}{4}\sin\phi\frac{\Gamma^2}{E_0^2+\Gamma^2}$, which reproduces the exact result for dc Josephson current through a resonant level. (3) The current formula is for the gauge choice of $V_L=V_N$ and $V_R=0$. It is easy to see that the formula is invariant under the transformation $V_N\to -V_N$, $N\to -N$, and $E_0\to E_0+V_N$, which corresponds to the gauge choice $V_L=0$ and $V_R=-V_N$. (4) $E_m^{\pm}$ are the poles of $G^r(\epsilon)$, which can be interpreted as PAAS. Obviously, E_m^+ and E_m^- contribute to the supercurrent with opposite signs. One can expect that resonant structures will appear near $E_0=\frac{1}{2}m\omega$. (5) Bessel functions enter the current formula not in the simple square form $J_n^2(x)$. It can be shown that $\bar{I}(\alpha_L,\alpha_R,\phi)=\bar{I}(-\alpha_L,-\alpha_R,\phi)$ for N even and $\bar{I}(\alpha_L,\alpha_R,\phi)=\bar{I}(-\alpha_L,-\alpha_R,\phi+\pi)$ for N odd.

Numerical Results

Firstly, we discuss the case where MW is applied symmetrically to the left and right tunnel barriers, i.e., $\alpha_L=\alpha_R$. Fig. 2 shows the curves of $I_c\equiv\bar{I}(\phi=\pi/2)$ versus $E_0-\frac{V_N}{2}$ at bias voltage $V_N=N\omega/2$ (N from 0 to 4), with symmetric MW strengths $\alpha_L=\alpha_R$. Parameters are: $\omega=1$, $\Gamma_L=\Gamma_R=0.02$, $k_BT=0.001$. $V_L=V_N$ and $V_R=0$ are set as a convention. Four features are noteworthy in the plot: (1) The curve is symmetric (anti-symmetric) with respect to $E_0=\frac{V_L+V_R}{2}=\frac{V_N}{2}$ for N even (odd). (2) Photon-assisted resonant structures appear near $E_0=V_L-\frac{1}{2}m\omega=V_R+\frac{1}{2}m\omega$. (3) There are two types of structures: single peak (dip) and peak-dip pair. (4) The structures grow with MW strength non-monotonically.

Feature (1) is due to the relation $I(E_0,\phi)=I(V_N-E_0,\phi)$ for N even and $I(E_0,\phi)=-I(V_N-E_0,\phi)$ for N odd. Feature (2) can be understood in terms of PAAS shown in Figs. 1b and 1c: electrons and holes are Andreev reflected back and forth by the superconducting gaps. When E_0 is near the position $V_L-\frac{1}{2}m\omega=V_R+\frac{1}{2}m\omega$, quasiparticles may exchange m photons with the MW field at the left tunnel barrier and exchange m photons at the right, so that the trajectory is closed and bound states are formed. It is these PAAS that play the role of carrying supercurrent. Understanding features (3) and (4) requires quantitative analysis of each photon-assisted structure. To proceed, we decompose the total current into $\bar{I}=\sum_m I_m$. Let $E_0=\frac{1}{2}m\omega+\delta$, and expand I_m near $E_0=\frac{1}{2}m\omega$, one can obtain

$$I_m=\frac{2e\Delta^2}{4}\sin\phi\frac{J_{m-N}(2\alpha_L)J_m(2\alpha_R)}{|\Gamma_m|}\left[\frac{\delta}{\sqrt{\delta^2+|\Gamma_m|^2}}\pm\frac{\delta}{\sqrt{\delta^2+|\Gamma_m|^2}}\right]$$

The cancellation between I_m^+ and I_m^- results in two types of resonant structures: $C_1\frac{\delta}{\sqrt{\delta^2+|\Gamma_m|^2}}$ and $C_2\frac{\delta}{\sqrt{\delta^2+|\Gamma_m|^2}}$, corresponding to single peak (dip) and peak-dip pair. At zero temperature, C_1 and C_2 are proportional to

$$C_1 \propto \sum_l \left[\delta_{l,\frac{1}{2}m} J_l^2(\alpha_R) + \delta_{l,\frac{1}{2}(m-N)} J_l^2(\alpha_L)\right]$$

$$C_2 \propto \sum_{l>-\frac{1}{2}m} J_l^2(\alpha_R) - \sum_{l>\frac{1}{2}m} J_l^2(\alpha_R) + \sum_{l>-\frac{1}{2}(m-N)} J_l^2(\alpha_L) - \sum_{l>\frac{1}{2}(m-N)} J_l^2(\alpha_L)$$

The oscillatory nature of Bessel functions in these coefficients is responsible for feature (4). There is an interesting special case: $N = 2$, $m = 1$, $\alpha_L = \alpha_R$, for which $C_1 = C_2 = 0$ for arbitrary MW strengths. Correspondingly, the resonant structure near $E_0 - \frac{V_N}{2} = 0$ is always missing for $N = 2$.

Next, we investigate the case where MW is applied only to one of the tunnel barriers, i.e., $\alpha_L = 0$ and $\alpha_R \neq 0$. Fig. 3 shows the curves of $I_c \equiv \bar{I}(\phi = \pi/2)$ versus E_0 at bias voltage $V_N = N\omega/2$ (N from -4 to 4), with MW strengths $\alpha_L = 0$ and $\alpha_R = 1$. Other parameters are the same as in Fig. 2. The inset shows I_c at $E_0 = V_L$ versus the MW strength α_R. In contrast to the symmetric case, a single peak is pinned at $E_0 = V_L$. The reason is as follows: since no MW is applied to the left barrier, Andreev tunneling through this barrier occurs only when E_0 aligns with V_L, while photon-assisted processes are allowed at the right barrier irradiated by the MW field. The inset shows I_c versus MW strength α_R at $E_0 = V_L$ for $N > 0$ (the case of $N < 0$ can be easily deduced from $N > 0$). One can see in the plot that the peak height (including the sign) is predominately determined by the prefactor $J_{m-N}(2\alpha_L)J_m(2\alpha_R)$. For $\alpha_L = 0$, $J_{m-N}(2\alpha_L)$ requires $m = N$, and the peak height is proportional to $J_N(2\alpha_R)$. We note that there exist some regions of MW strength where I_c and V_N have opposite signs, which can be viewed as a quantum pump effect. However, this feature is dramatically different from the quantum pump effect in normal mesoscopic junctions.

Conclusion

To sum up, we have investigated the ac Josephson effect in resonant tunneling through mesoscopic superconducting junctions. We show that PAAS play an essential role in conducting supercurrent through a narrow resonance, in which the MAR trajectory is closed by exchanging photons with the MW field. On the Shapiro steps $2V = N\omega$, a dc component appears when the resonant level E_0 is near $V_L - \frac{1}{2}m\omega = V_R + \frac{1}{2}m\omega$, due to the formation of PAAS. In the limit $\Delta \to \infty$, the analytical result Eq. (26) is derived for the dc component of the Josephson current, which helps to understand new features of the ac Josephson effect in the case of resonant tunneling.

We have dropped Coulomb interaction in the calculation. The derived results are meaningful in the following senses: (1) They are directly applicable to systems where Coulomb blockade effects are negligible, i.e., $U \lesssim \Gamma \ll \Delta$ (U is the charging energy). This is possible by using a substrate with large dielectric constant to reduce U or using proper material as superconducting leads to obtain large Δ. (2) They are instructive for more complicated cases. We note that the concept of PAAS is also useful for the interacting case. For $U \to \infty$, the resonant level E_0 is effectively split into two resonances E_0 and $E_0 + U$. Similar to Figs. 1b and 1c, one can draw diagrams of closed MAR trajectories through these resonances. Moreover, Coulomb blockade can be partially removed by bias voltage or MW irradiation, as long as U is comparable to Δ. Obviously, calculation including interaction terms is much more difficult, and analytical results are almost impossible. Efforts along this line are still in progress.

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Figure Captions

Fig. 1 [Figure 1: see original paper] Schematic diagram of resonant tunneling through mesoscopic superconducting junctions. (a) Without MW irradiation, the MAR trajectory cannot be covered by a narrow resonance, and the tunnel current is exponentially small. (b) and (c) In the presence of MW irradiation, the MAR trajectory can be closed by emitting or absorbing photons. Two types of PAAS are formed, carrying supercurrent with opposite signs.

Fig. 2 [Figure 2: see original paper] The curves of $I_c \equiv \bar{I}(\phi = \pi/2)$ versus $E_0 - V_N/2$ at bias voltage $V_N = N\omega/2$ (N from 0 to 4), with symmetric MW strengths $\alpha_L = \alpha_R$. Parameters are: $\omega = 1$, $\Gamma_L = \Gamma_R = 0.02$, $k_BT = 0.001$. $V_L = V_N$ and $V_R = 0$ are set as a convention.

Fig. 3 [Figure 3: see original paper] The curves of $I_c \equiv \bar{I}(\phi = \pi/2)$ versus E_0 at bias voltage $V_N = N\omega/2$ (N from -4 to 4), with MW strengths $\alpha_L = 0$ and $\alpha_R = 1$. Other parameters are the same as in Fig. 2. The inset shows I_c at $E_0 = V_L$ versus the MW strength α_R. Diagrams for corresponding PAAS are also shown in the plot.

Note: Figure translations are in progress. See original paper for figures.

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