Title

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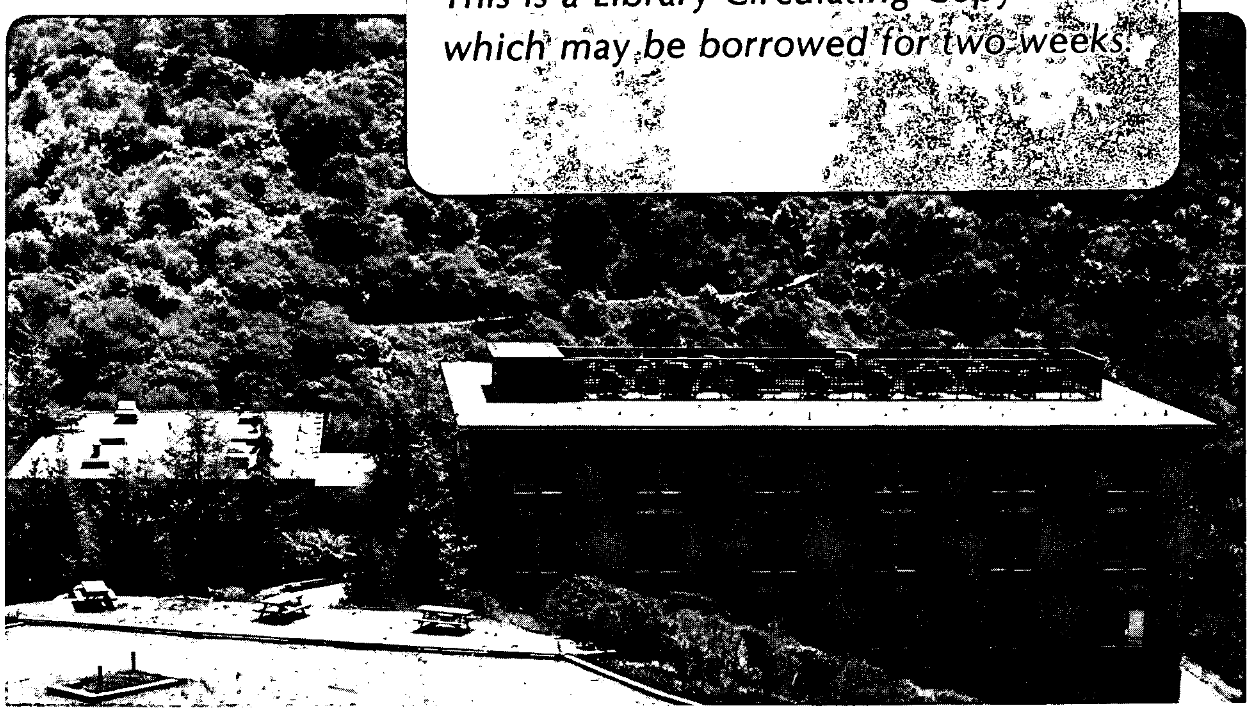

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DISCLAIMER

EXPERIMENTAL OBSERVATION OF THE QUANTUM BEHAVIOR OF A MACROSCOPIC DEGREE OF FREEDOM

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Contents

1. Introduction

It is usually taken for granted that a macroscopic degree of freedom, such as the position of the center of gravity of a billiard ball, never displays quantum effects. However, as Leggett shows theoretically in this volume [1], this needs not always be so. We performed at Berkeley a serie of experiments [2] that, we believe, demonstrates the quantum behavior of one macroscopic degree of freedom, namely the phase difference across a current biased Josephson junction. In his review, Leggett explains why quantum phenomena affecting a macroscopic variable are possible in circuits involving Josephson junctions.

Here, we will focus on the practicalities involved in demonstrating that a macroscopic variable indeed behaves quantum-mechanically. Since the details of the experiments have been described elsewhere [2], we will try here only to sketch their general principles.

First, what do we mean when we say we observe experimentally the quantum behavior of a system? For a microscopic system like an atom the situation is simple. The properties of each particle (electrons and nucleus) are known independently and the dynamical equations of the system can be obtained from first principles. To make predictions one treats these equations classically or quantum mechanically. Of course, for the atom, it is the quantum predictions, not the classical ones that correspond to the observed properties. As for a macroscopic system, whose complete description can never be obtained in every detail, the situation is more complicated and involves a two-step procedure. In the first step one makes measurements in a regime where thermal fluctuations are large compared with the expected quantum fluctuations. One gets thereby the parameters entering in the classical equations of motion that characterize completely the system. In the second step, the thermal fluctuations are suppressed. The new properties are then measured. Finally one compares the results obtained in the second step with the predictions of the quantum theory based on the results of the first step. If they always agree, the quantum behavior of the macroscopic system is said to have been observed.

Two obvious pitfalls must be avoided in this procedure. First, one must make sure that all the classical parameters involved later in the quantum predictions are measured correctly. Second, one must verify, while the measurement is being made, that all thermal and parasitic fluctuations that could mimic quantum fluctuations are indeed suppressed.

In pioneering experiments [3,4,5,6] on macroscopic quantum tunneling involving Josephson junctions great care was taken in minimizing uncontrolled perturbations that could destroy any quantum effect. The results of these experiments were consistent with a quantum mechanical interpretation of the behavior of the phase difference of the junctions. Recently, more sophisticated experiments investigated the question of the influence of dissipation on tunneling [7,8] which had been calculated theoretically

[1, 9].

However, a persistent difficulty in all these experiments has been the correct determination of the parameters of the principal component of the macroscopic system, namely the Josephson junction. Without a precise knowledge of these parameters no quantitative comparison with quantum theories can be made. Some of the important parameters were treated as ajustable parameters in the predictions or were estimated from measurements of properties not directly connected to the dynamical properties involved in the tunelling experiment. This is why we decided to devise new experimental techniques to ensure that all the parameters of the junction would be measured in situ, in particular under those conditions that guarantee that no spurious noise reach the junction.

In the following the emphasis will be put on the particular procedures we used to solve the two problems of noise shielding and parameter determination. Before discussing these central issues, we begin with a short description of the macroscopic system we investigated, the current biased Josephson junction.

2. Classical dynamics of the current biased Josephson junction

In experiments on Josephson junctions biased with a constant current such as ours, the macroscopic degree of freedom is the phase difference across the junction. A Josephson junction is made with two superconductors separated by a very thin oxide layer. For the purpose of understanding our experiment, the underlying physics of the superconductors and their coupling through the oxide barrier is unimportant and the junction can be safely thought of as an ideal <u>non-linear</u> inductor in parallel with a capacitor (see Fig. 1). An ideal linear inductor would be a simple coil made of superconducting wire and would be characterized by an energy quadratic in the flux through the coil. Instead, the ideal nonlinear inductor that constitutes a Josephson junction has an energy given by

$$E=-I_0(\Phi_0/2\pi)\cos(\delta) \qquad (1)$$

where Φ_0 is the flux quantum. The variable δ is the phase difference across the junction. The quantity $\varphi=(\Phi_0/2\pi)\delta$ is analogous to the flux through the inductor in the sense that the current and voltage for this element satisfy $i=dE/d\varphi$ and $v=d\varphi/dt$.

The junction by itself is thus characterized by two parameters : the critical current I_0 and the capacitance C. We will see that the coupling of the junction to the measuring apparatus involves additional parameters.

When it is biased with a constant current I, the junction is equivalent to the model of a particle moving in a tilted washboard potential (see Fig. 2). The mass of the particle is the capacitance C. The tilt of the washboard is the ratio I/I_0. When

$I<I_0$ the potential has relative minima and the particle can be in two states. In the so-called zero voltage state the particle is confined in one well (see Fig. 2a). The average velocity and hence the dc voltage across the junction is zero. In the other state called the voltage state, the particle runs down the washboard at an average constant velocity determined by frictional forces (see Fig. 2b). In our experiment the velocity is solely limited by the breaking of Cooper pairs in the superconductors and correspond to the gap voltage which is of the order of 2 mV.

The zero voltage state is metastable and eventually decays into the voltage state. It is easy to determine the decay rate of the zero voltage state by monitoring the voltage across the junction. The experiment occurs as follows : Initially $I=0$ and the particle representing the junction phase difference is in a relative minimum of the washboard potential. We ramp the bias current from 0 to $I<I_0$, and wait until a voltage appears across the junction. This voltage is the signature that the particle has escaped from the well and accelerated down the washboard potential. It is worth noting that the acceleration of the particle "amplifies" the very small voltage pulse associated with the escape event. This built-in amplification process is a particularly attractive feature of the current biased junction.

Because the acceleration process is so fast, the time at which voltage was measured across the juntion can be taken as the time at which the particle escaped. We repeat the experiment a large number of times to get the average escape rate.

By measuring the escape rate as a function of the bias current which controls the shape and the size of the potential well one gets access to some aspects of the behavior of the particle inside the well. To get a complete picture of what is happening to the particle inside the well one needs extra "knobs" on the experiment. This is where our experiments fundamentally differs from all the others : we add to the dc current bias a weakly perturbing ac current source and measure the <u>change</u> in the escape rate induced by this perturbation. The variations of this change with the frequency and the bias current provide additional information that, when combined with the first measurement, enables one to reconstruct the <u>dynamics</u> of the particle in the well.

We have supposed up to now that the current source and the voltmeter are ideal. In practice they have always some finite impedance in parallel with them, in particular the impedance due to the electromagnetic coupling between the two leads connecting the instruments to the junction. Even though these impedances may be negligible at low frequencies and may not affect a quasi-static measurement such as an I-V characteristic, they become of prime importance at the frequencies relevant for the motion of the particle in the well which are in the microwave range. The junction will thus see an effective admittance in parallel with it. As we will see in the next section, it is possible to arrange the circuit so that this admittance behaves essentially like a resistor in parallel with a capacitor. This picture is valid in a broad frequency range around the characteristic frequency of the system. The effect of the external shunting capacitance is simply to renormalize the capacitance of the junction. The effect of the resistance is two-fold. One effect is to damp the motion of the particle in the well. The other is to act as a thermal bath and to induce fluctuations in the motion of the particle. Two additional parameters must then be introduced to take into account the coupling of the junction with the measuring apparatus : the resistor R and its temperature T. The temperature of the junction itself is unimportant as long as its intrinsic damping is negligible when compared with the damping provided by the external resistor R - this is always the case in our experiment.

The above description of our current biased Josephson junction corresponds to the circuit of Fig. 3 and leads to the following classical equation of motion :

$$C\,(\Phi_0/2\pi)^2\,\ddot{\delta} + R^{-1}(\Phi_0/2\pi)^2\,\dot{\delta} + \frac{\partial}{\partial\delta}U(\delta) = (\Phi_0/2\pi)I_N(t) \qquad (2)$$

where $U(\delta)$ is the tilted cosine potential

$$U(\delta) = (-I_0\Phi_0/2\pi)\left[\cos\delta + (I/I_0)\,\delta\right] \qquad (3)$$

and where the current noise I_N satisfies

$$\int_{-\infty}^{+\infty} \langle I_N(t)\,I_N(0)\rangle_T\, e^{i2\pi\nu t}\,dt = \frac{2k_BT}{R} \qquad (4)$$

In practice the bias current I is very close to the critical current and the potential from which the particle escapes is very nearly cubic (see Fig. 4). The barrier height ΔU and the oscillation frequency ω_p at the bottom of the well (plasma frequency) are two useful independent parameters that describe completely the potential and are given, to a very good approximation, by :

$$\Delta U = (4\sqrt{2}\, I_o \Phi_o / 6\pi)(1 - I/I_o)^{3/2} \tag{5}$$

and

$$\omega_p = (\sqrt{2}\, 2\pi I_o / \Phi_o C)^{1/2} (1 - I/I_o)^{1/4} \tag{6}$$

The damping due to the resistor is conveniently described by the dimensionless quality factor of the small oscillations at the bottom of the well

$$Q = RC\omega_p \tag{7}$$

Before turning to the problem of determining the parameters I_o, C, R and T, we will describe some crucial features of our apparatus, namely the filters that isolate the junction from unwanted perturbation coming from the measurement circuitry.

3. Decoupling the junction from ambient noise

It goes without saying that the junction must be completely shielded from parasitic electromagnetic sources like radio stations. But this is far from being sufficient. Thermal current noise from the measuring apparatus at room temperature - i.e. the current source and the voltage amplifier - must also be prevented from reaching the junction through its leads. This is achieved by interposing a series of filters between the junction and the measuring apparatus. These filters screen out all frequencies except a very narrow range around zero so that we are nevertheless able to vary the bias current and detect the voltage rise due to the junction switching to the voltage state. The range is chosen to be sufficiently narrow so that noise in that frequency band is negligible and sufficiently wide that a good time resolution of the lifetime of the zero voltage state is achieved.

In practice such a low pass wideband filter can only be made conveniently with filter elements having substantial dissipation. Because of this dissipation, the filter will itself produce noise and each filter in the chain must be at a lower temperature than the preceding one. Because of its importance we developed a special type of filter for the experiment. It consists of a spiral coil of wire inside a copper tube filled with copper powder with a grain size of about 30 μm. Since each grain is insulated from its neighbour by an oxide layer, the effective area of the copper is enormous and thus provides substantial skin effect losses even at the lowest temperatures. The lowest part of the radio-frequency spectrum was attenuated with a classical RC network. The chain of filters provided a total attenuation of more than 200 dB.

The last filter was carefully engineered since it imposes the damping and the fluctuations of the phase difference of the junction. It consists of an attenuating coaxial line that also relies on copper powder for the source of damping. The copper powder was thermalized carefully by injecting epoxy between the grains. The junction is mounted as closely as possible to the end of the line to ensure that the impedance discontinuity between the junction and the line occurs on a distance small compared with the wavelength at the plasma frequency. Thus, the impedance seen by the junction behaved essentially like a parallel RC combination with no important spurious resonances.

4. Determination of junction parameters using classical phenomena

4.1. Capacitance and resistance

The parameters ω_p and Q were determined by resonant activation. This phenomenon involves the enhancement of the escape rate by a microwave irradiation of the junction. This enhancement is determined by taking the ratio of the escape rate measured in the presence of microwave irradiation of power P and of the escape rate measured in the same conditions but without microwaves. Below the plasma frequency, the enhancement is a smoothly increasing function of the irradiation frequency. But when a frequency just slightly below the plasma frequency is reached, the enhancement drops steeply to zero. As numerical simulations [10,11] and analytical calculations [12,13,14] have shown, the plasma frequency is given, apart from a small correction, by that frequency at which the enhancement drops steeply while Q is given by the range of frequency over which the drop occurs. Experimentally, it is easier to keep the irradiation frequency fixed and to vary the plasma frequency by varying the bias current (see Eq. 6). An example is shown in Fig. 5. At the resonant current I_{res}, the plasma frequency is equal to the applied current. From Eq. 6 we can deduce the plasma frequency at the other values of the bias current. The value of Q can be inferred from the width in current of the roll-off of the enhancement.

4.2. Critical current and temperature

In the thermal regime the escape of the particle from the well occurs by thermal activation. The rate is given by [15]

$$\Gamma_t = a_t(\omega_p/2\pi)\exp(-\Delta U/k_BT), \tag{8}$$

where the prefactor a_t is given by

$$a_t = 4/[(1 + Qk_BT/1.8\,\Delta U)^{1/2} + 1]^2. \tag{9}$$

We determine I_0 and T from the dependence of the escape rate on the bias current. As is evident from Eq. 5 and Eq. 8 a plot of the experimentally determined quantity $\{\ln[\omega_p(I)/2\pi\Gamma(I)]\}^{2/3}$ vs. I should, neglecting departures of a_t from unity be a straight line with slope scaling as $T^{2/3}$ that intersects the current axis at I_0. Figure 6 shows three examples of such plots. As expected lines drawn through the data (dashed lines) intersect the current axis at very nearly the same point. Taking into account the correction due to the departure of a_t from unity we obtain a temperature independent value of the critical current (arrow on Fig. 6). We find excellent agreement between the temperature measured from the current dependence of the escape rate and the temperature measured by our thermometers. The fact that the data fall indeed on a straight line is also an additional check that the escape occurs through thermal activation at a well defined temperature.

We have described how we obtain all the system parameters using the properties of the escape in the classical regime. We now go to the quantum regime by reducing the temperature and remeasure the same properties.

5. Measurements in the quantum regime

5.1. Quantized energy levels

At low temperatures, one has to decrease substantially the barrier height to get a short enough lifetime of the zero voltage state. Thus, there are only several levels in the well. For temperatures T slightly greater than $h\omega_p/2\pi k_B$ [1,9], the escape in the absence of microwaves occurs essentially via thermal activation through these discrete states to the continuum of states above the potential barrier. In the presence of microwave power P however, the population of the excited states at the top of the well increases. The particle then escapes with a rate $\Gamma(P)$ which is greater than $\Gamma(0)$. Consequenty, one expects a resonant enhancement of the escape rate when the irradiation frequency matches a transition frequency between two energy levels. As in the classical regime, we keep the frequency fixed and vary the bias current I, thereby changing continuously all the level spacings. One can in this way determine spectroscopically the position of the energy levels in the well. Our results for a C= 47 pF junction are plotted in Fig. 7a.

These results can be compared with a quantum calculation of the energy levels based on an in situ measurement of the parameters ΔU and ω_p. An unambiguous agreement between the experimental value of the relative positions of the resonances and their theoretical predictions is found (see Fig. 7b). The relatively small shift in the absolute value of the position of the resonances is due to the uncertainty in the determination of the critical current (the resulting uncertainty in the 0 -> 1 transition frequency is shown by the dotted line).

In this experiment, the junction is only slightly damped with a Q factor around 80. The relative width of the 0 -> 1 resonance in the quantum regime is approximatively equal to 1/Q, as predicted by theory [16].

5.1. Quantum tunneling

In the absence of microwaves and for temperatures below $h\omega_p/2\pi K_B$, escape from the well occurs via quantum tunneling through the potential barrier. The predicted quantum tunneling rate is [1] :

$$\Gamma_q = a_q \frac{\omega_p}{2\pi} \exp\left[- 7.2 \frac{\Delta U}{\hbar\omega_p} \left(1 + \frac{0.87}{Q} + \ldots\right)\right], \tag{10}$$

where

$$a_q = [120\pi(7.2\ \Delta U/\hbar\omega_p)]^{1/2}. \tag{11}$$

Eq. 10 is written so that a direct comparison with Eq. 8 can be made. The usual WKB result is obtained by setting Q -> I in

Eq. 10.

Our results are shown in Fig. 8 where the escape rate is plotted as an escape temperature T_{esc} (defined through the relation $\Gamma = (\omega_p/2\pi)\exp(-\Delta U/KT_{esc})$ versus temperature. This measurement was done for a C = 6 pF junction which could be placed more firmly in the quantum regime (at the expense of a somewhat greater damping : Q = 30 ; however the damping was not strong enough to affect significantly the escape rate). The white dots correspond to a measurement done with a magnetic field on the junction which has the effect of reducing its critical current. With a reduced critical current, the plasma frequency is lower and the junction should behave classically down to the lowest temperatures. Indeed the escape rate followed the classical prediction (solid line) showing that no significant amount of external noise was reaching the junction. When the critical current is restored to its original value, we measured a higher value of the escape rate (solid dots). The escape rate became temperature independent at low temperatures which is expected if the escape mechanism is dominated by quantum fluctuations. This interpretation is supported by the quantitative agreement between this limiting value of the measured escape rate and its theoretical prediction (arrow shows the prediction of Eq. 10; we have also indicated the error bar of the prediction due to the uncertainties in the classical measurements of the parameters). We note that in our experiment the predicted effect of dissipation on tunneling is too small to be discernible.

6. Conclusion

Our measurements of the lifetime of the zero-voltage state of a current biased Josephson junction agree quantitatively with predictions based on quantum theory, all the relevant parameters being measured in situ. This shows that the phase difference across a Josephson junction behaves as a macroscopic quantum variable.

Can one conclude that macroscopic degrees of freedom always obey Quantum Mechanics? At this point, the experimentalist can choose between two attitudes: the Utopist's or the Pragmatist's.

For the Utopist, who worries about the "crazy framework of (probability) amplitudes" [17] of Quantum Mechanics, the answer to the above question is "no". He must pursue the search for an experiment that would show the limits of validity of Quantum Mechanics. He is led to test, for example, the existence of macroscopic quantum coherence [1] or the superposition principle involving highly excited states [18].

For the Pragmatist, who wants to use this new macroscopic Quantum Mechanics, the answer to the question is "yes, probably". He is inclined to explore new superconducting circuits that would perform quantum signal processing such as squeezing [19] and perhaps implement the so-called back action evasion amplifiers that are needed for the detection of gravity waves [20]. Finally, the Pragmatist can dream of building exotic macroscopic "atoms with wires" that would display new quantum phenomena with no equivalents in the microscopic world.

Acknowledgments

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Figure captions

Fig. 1 : Comparison between a Josephson junction, which consists of a non-linear inductor in parallel with a capacitor, and a linear LC circuit.

Fig. 2 : The two states of a Josephson junction biased with a current $I < I_0$; in a), $\langle \dot{\delta} \rangle = 0$ (zero-voltage state) while in b), $\langle \dot{\delta} \rangle \neq 0$ (finite voltage state)

Fig. 3 : Equivalent circuit for our junction coupled to its measurement apparatus which has a finite impedance at microwave frequencies. The resistor R models the real part of this impedance. The capacitor C includes a contribution from the imaginary part of the impedance.

Fig. 4 : Cubic potential from which the particle representing the junction phase difference escapes.

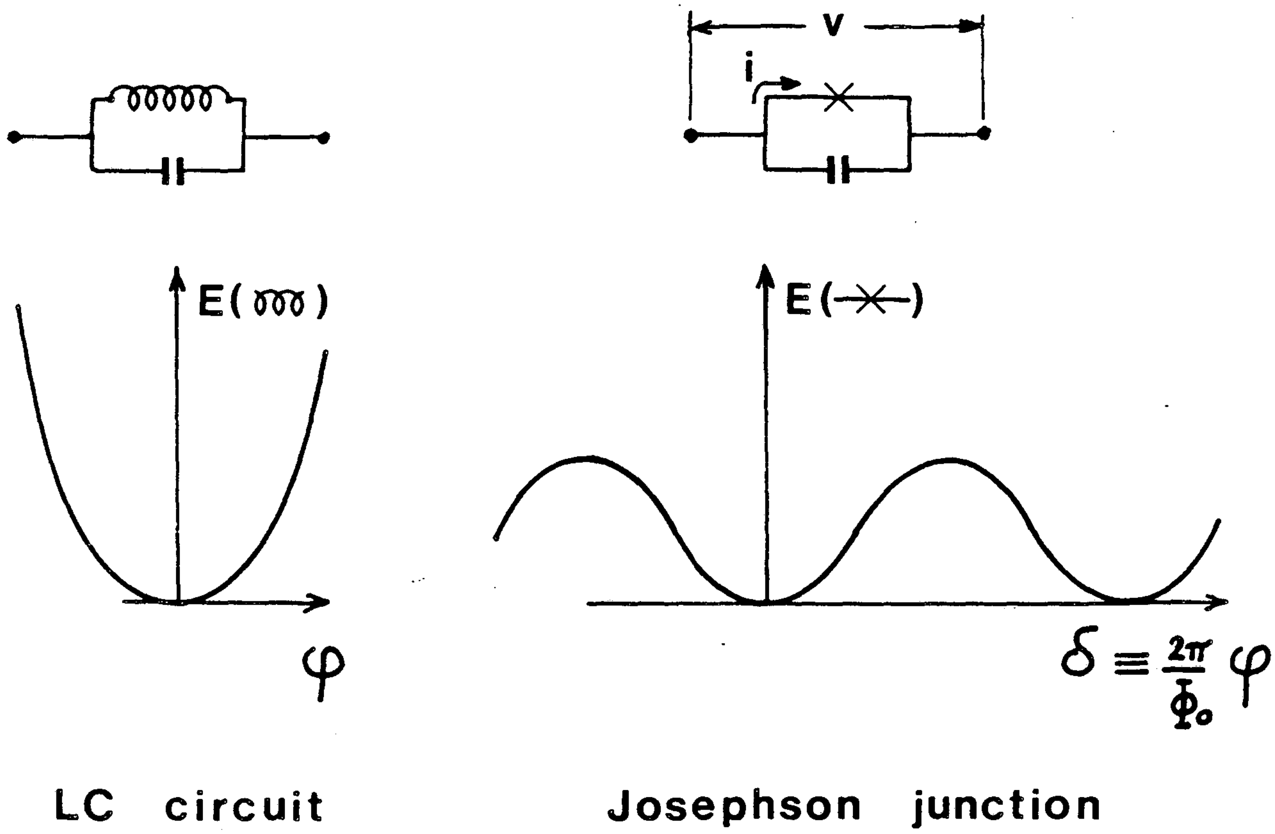
Fig. 5 : Microwave induced enhancement of the escape rate measured as a function of the bias current. $\Omega/2\pi$ is the irradiation frequency.

Fig. 6 : Plot of the 2/3 power of the logarithm of the escape rate vs. bias current. On that plot, data should fall on a straight line. Dashed lines are least-squares fit of the data.

The arrow indicates the position of the critical current.

Fig. 7 : a) Microwave induced enhancement of the escape rate as function of bias current for $k_BT/\hbar\omega_p = 0.29$. Arrows indicate positions of resonances. Inset represents the corresponding transitions between energy levels. b) Calculated energy level spacings. Dotted lines indicate uncertainties in the E_{01} curve due to errors in the determination of junction parameters. Arrows indicate values of bias current at which resonances are predicted.

Fig. 8 : Escape temperature plotted against temperature for two values of the critical current. The solid line is the classical prediction $T_{esc} = 0.95T$ (Eq. 8). The arrow is the quantum prediction (Eq. 10).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i
E( )
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$\delta \equiv \frac{2\pi}{\Phi_0}\varphi$
LC circuit
Josephson junction

Fig. 1

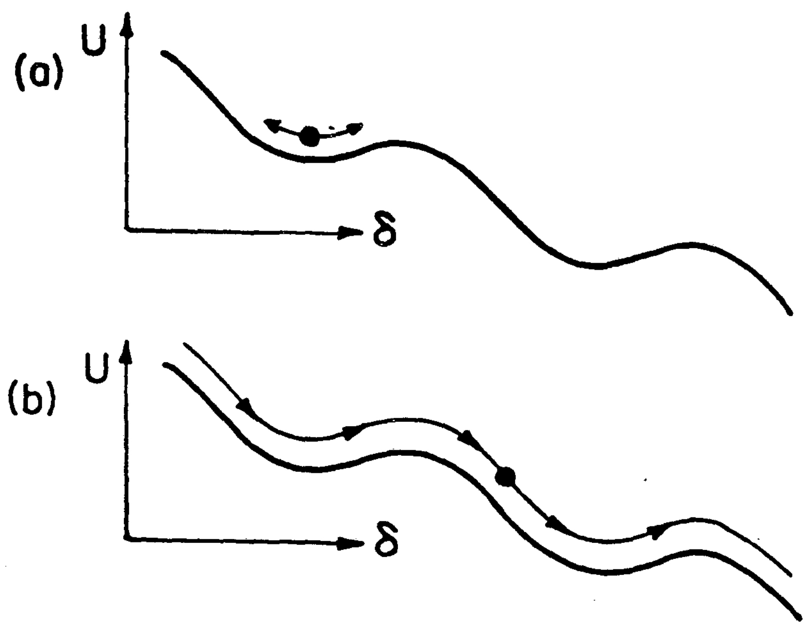


Fig. 2

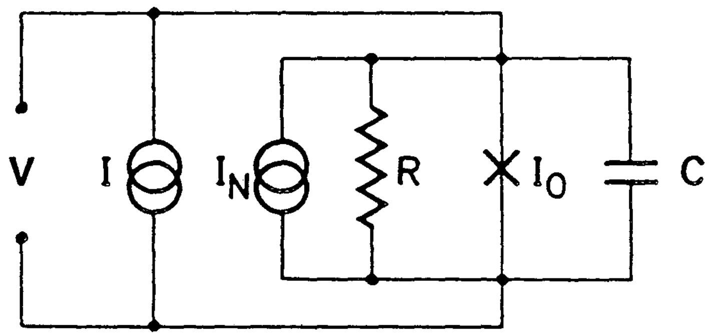


Fig. 3

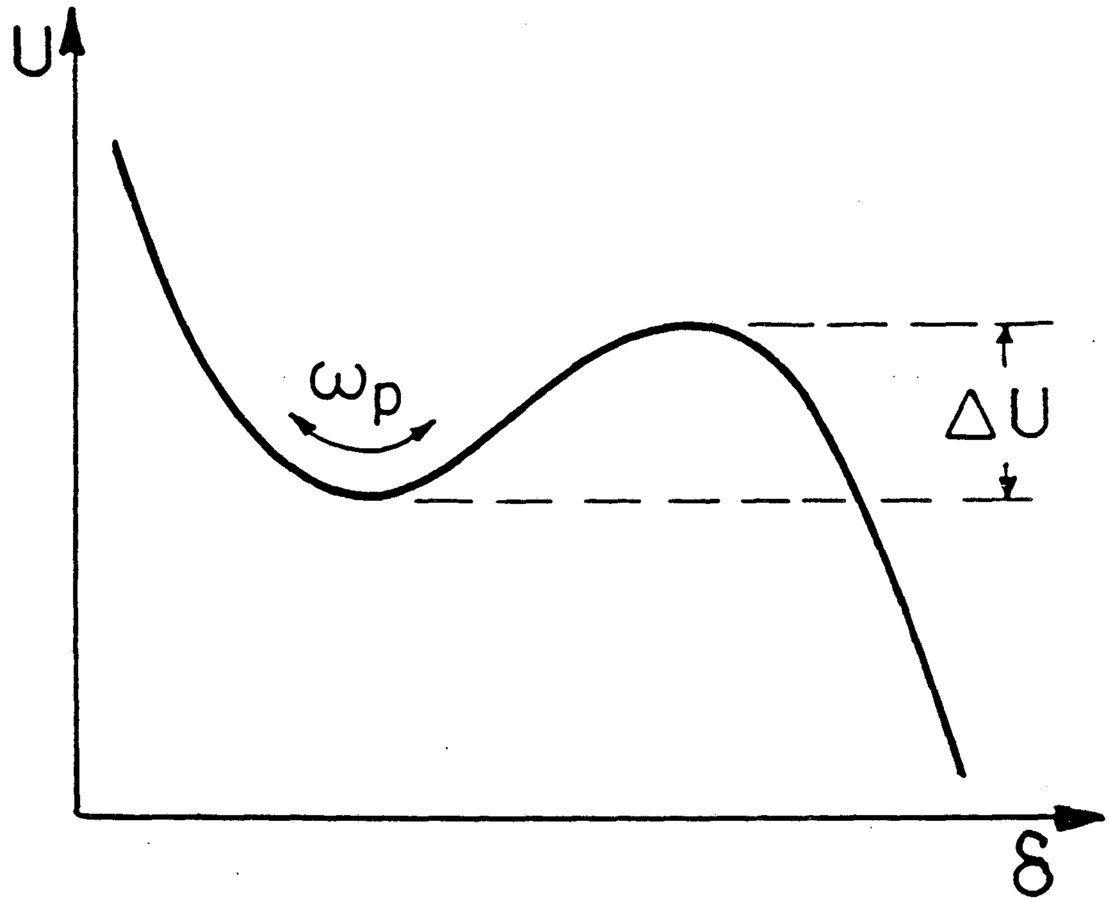


Fig. 4

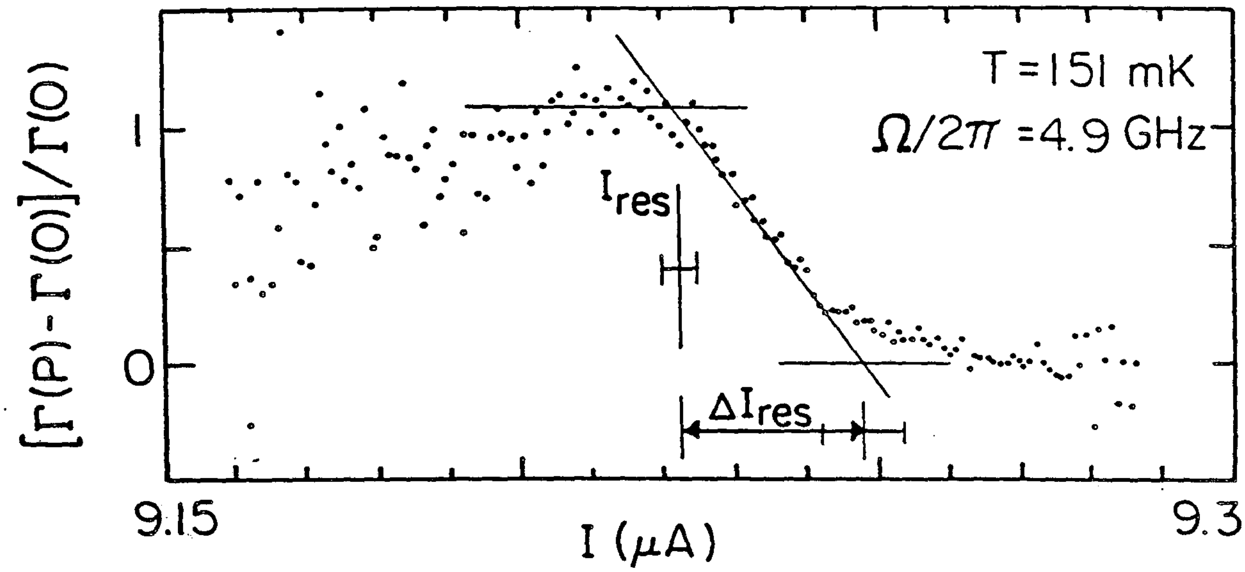


Fig. 5

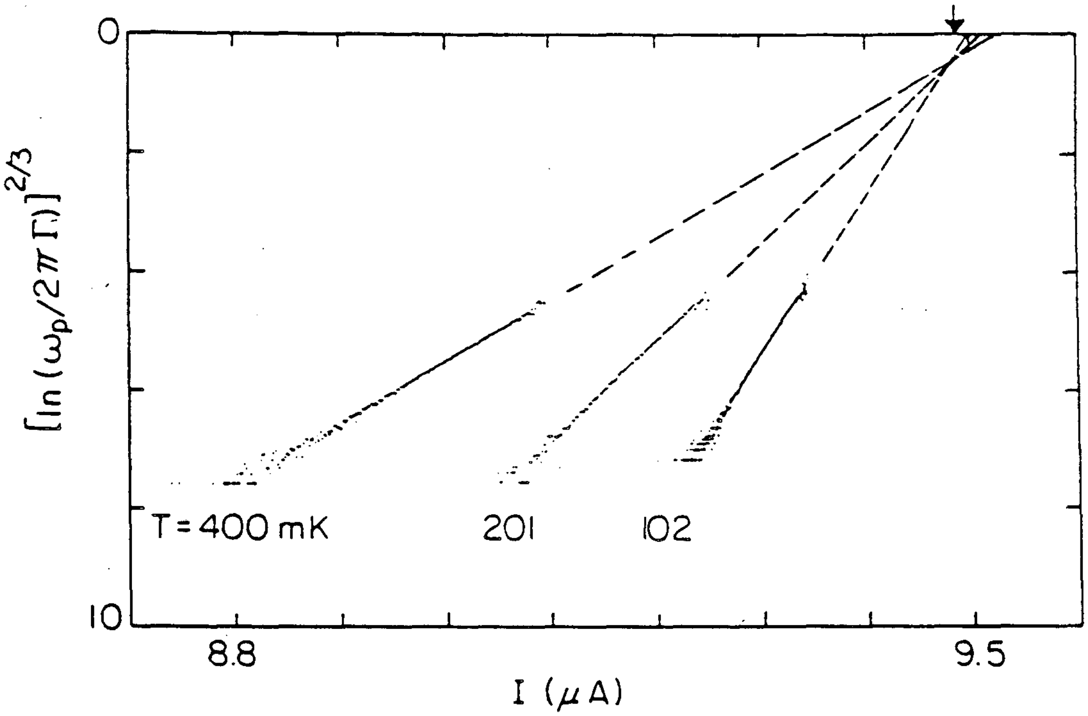


Fig. 6

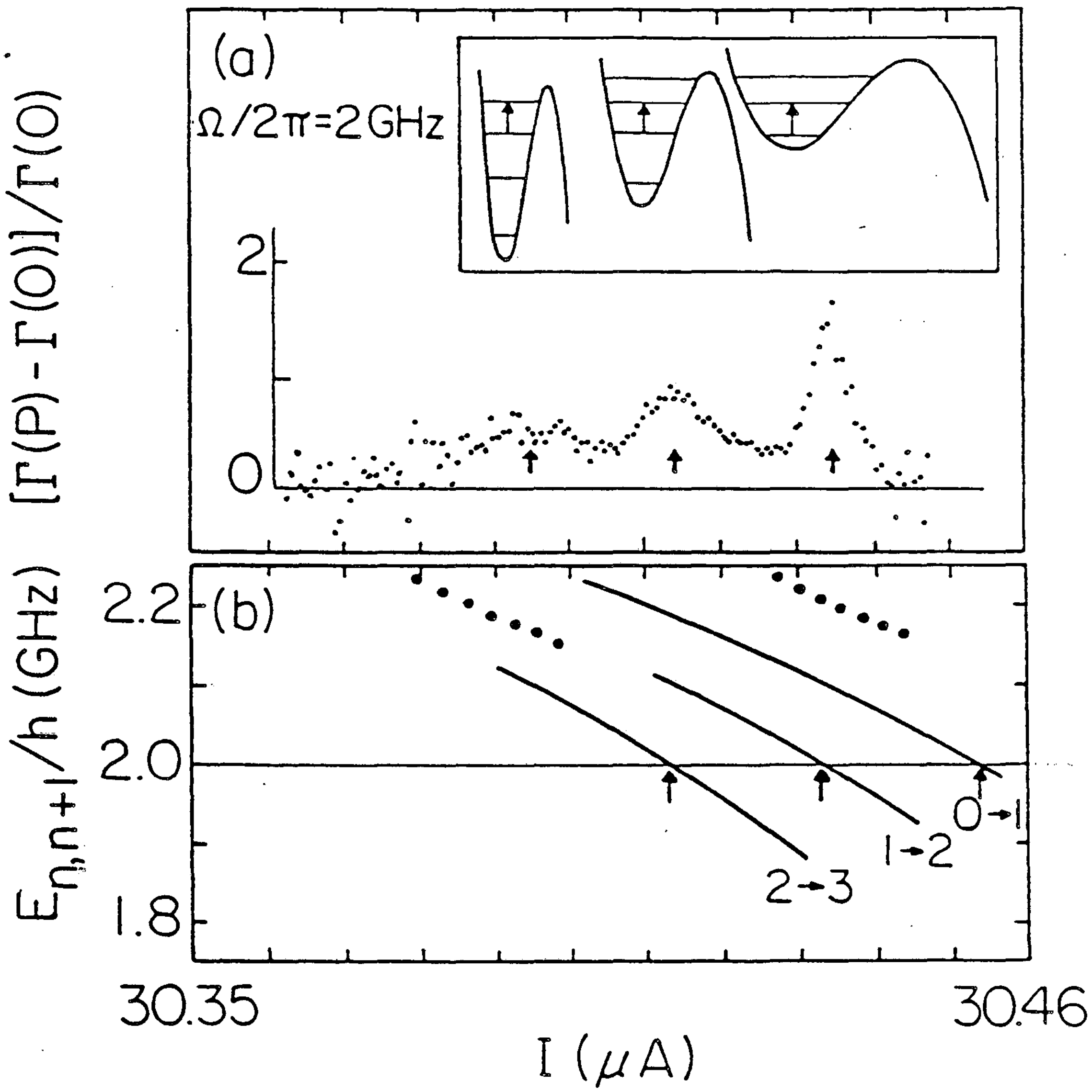
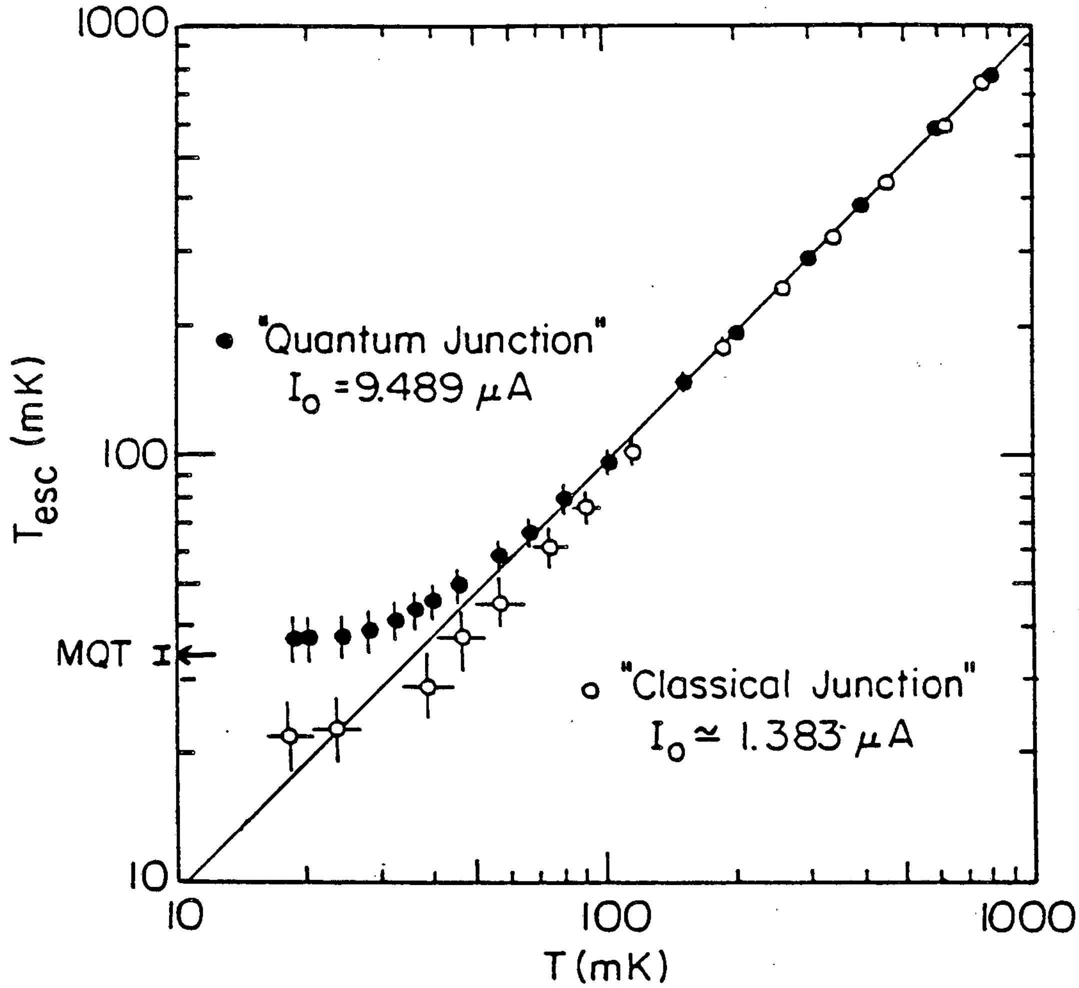


Fig. 7



1000
100
10
10
100
1000
T (mK)
T_{esc} (mK)
MQT
"Quantum Junction"
I_0 =9.489 μA
"Classical Junction"
$I_0 \simeq$ 1.383 μA

Fig. 8